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# PROPERTIES OF HADRONS IN THE NUCLEAR MEDIUM

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## Abstract

This review is devoted to the discussion of hadron properties in the nuclear medium and its relation to the partial restoration of chiral symmetry. Special attention is given to disentangle in-medium effects due to conventional many-body interactions from those due to the change of the chiral condensate. In particular, we shall discuss medium effects on the Goldstone bosons (pion, kaon and eta), the vector mesons (rho, omega, phi), and the nucleon. Also, for each proposed in-medium effect the experimental consequence and results will be reviewed.

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## 1 INTRODUCTION

The atomic nucleus provides a unique laboratory to study the long range and bulk properties of QCD. Whereas QCD is well tested in the perturbative regime rather little is known about its properties in the long range, nonperturbative region. One of the central nonperturbative properties of QCD is the spontaneous breaking chiral symmetry in the ground state resulting in a nonvanishing scalar quark condensate,  $\langle \bar{q}q \rangle \neq 0$ . It is believed and supported by lattice QCD calculations [1] that at temperatures around 150 MeV, QCD undergoes a phase transition to a chirally restored phase, characterized by the vanishing of the order parameter, the chiral condensate  $\langle \bar{q}q \rangle$ . This is supported by results obtained within chiral perturbation theory [2, 3]. Effective chiral models predict that a similar transition also takes place at finite nuclear density.

The only way to create macroscopic, strongly interacting systems at finite temperature and/or density in the laboratory is by colliding heavy nuclei at high energies. Experiments carried out at various bombarding energies, ranging from 1 AGeV (BEVALAC, SIS) to 200 AGeV (SPS), have established that one can generate systems of large density but moderate temperatures (SIS, BEVALAC), systems of both large density and

temperature (AGS) as well as systems of low density and high temperatures (SPS). Therefore, a large region of the QCD phase diagram can be investigated through the variation of the bombarding energy. But in addition the atomic nucleus itself represents a system at zero temperature and finite density. At nuclear density the quark condensate is estimated to be reduced by about 30% [2, 3, 4, 5] so that effects due to the change of the chiral order parameter may be measurable in reactions induced by a pion, proton or photon on the nucleus.

Calculations within the instanton liquid model [6] as well as results from phenomenological models for hadrons [7] suggest that the properties of the light hadrons, such as masses and couplings, are controlled by chiral symmetry and its spontaneous as well as explicit breaking. Confinement seems to play a lesser role. If this is the correct picture of the low energy excitation of QCD, hadronic properties should depend on the value of the chiral condensate  $\langle \bar{q}q \rangle$ . Consequently, we should expect that the properties of hadrons change considerably in the nuclear environment, where the chiral condensate is reduced. Indeed, based on the restoration of scale invariance of QCD, Brown and Rho have argued that masses of nonstrange hadrons would scale with the quark condensate and thus decrease in the nuclear medium [8]. This has since stimulated extensive theoretical and experimental studies on hadron in-medium properties.

By studying medium effects on hadronic properties one can directly test our understanding of those non-perturbative aspects of QCD, which are responsible for the light hadronic states. The best way to investigate the change of hadronic properties in experiment is to study the production of particles, preferably photons and dileptons, as they are not affected by final-state interactions. Furthermore, since vector mesons decay directly into dileptons, a change of their mass can be seen directly in the dilepton invariant mass spectrum. In addition, as we shall discuss, the measurement of subthreshold particle production such as kaons [9] and antiprotons [10] may also reveal some rather interesting in-medium effects.

Of course a nucleus or a hadronic system created in relativistic heavy ion collisions are strongly interacting. Therefore, many-body excitations can carry the same quantum numbers as the hadrons under consideration and thus can mix with the hadronic states. In addition, in the nuclear environment ‘simple’ many-body effects such as the Pauli principle are at work, which, as we shall discuss, lead to considerable modifications of hadronic properties in some cases. How these effects are related to the partial restoration of chiral symmetry is a new and unsolved ques-

tion in nuclear many-body physics. One example is the effective mass of a nucleon in the medium, which was introduced long time ago [11, 12] to model the momentum dependence of the nuclear force, which is due to its finite range, or to model its energy dependence, which results from higher order (2-particle - 1-hole etc.) corrections to the nucleon self energy. On the other hand, in the relativistic mean-field description one also arrives at a reduced effective mass of the nucleon. According to [13], this is due to so-called virtual pair corrections and may be related to chiral symmetry restoration as suggested by recent studies [3, 14, 15, 16].

Another environment, which is somewhat ‘cleaner’ from the theorists point of view, is a system at finite temperature and vanishing baryon chemical potential. At low temperatures such a system can be systematically explored within the framework of chiral perturbation theory, and essentially model-independent statements about the effects of chiral restoration on hadronic masses and couplings may thus be given. In the high temperature regime eventually lattice QCD calculations should be able to tell us about the properties of hadrons close to the chiral transition temperature. Unfortunately, such a system is very difficult to create in the laboratory.

Since relativistic heavy ion collisions are very complex processes, one has to resort to careful modeling in order to extract the desired in-medium correction from the available experimental data. Computationally the description of these collisions is best carried out in the framework of transport theory [17, 18, 19, 20, 21, 22, 23] since nonequilibrium effects have been found to be important at least at energies up to 2 AGeV. At higher energies there seems to be a chance that one can separate the reaction in an initial hard scattering phase and a final equilibrated phase. It appears that most observables are dominated by the second stage of the reaction so that one can work with the assumption of local thermal equilibrium. However, as compared to hydrodynamical calculations, also here the transport approach has some advantages since the freeze out conditions are determined from the calculation and are not needed as input parameter. As far as a photon or a proton induced reaction on nuclei is concerned, one may resort to standard reaction theory. But recent calculations seem to indicate that these reaction can also be successfully treated in the transport approach [24, 25, 26]. Thus it appears possible that one can explore the entire range of experiments within one and the same theoretical framework which of course has the advantage that one reduces the ambiguities of the model to a large extent by exploring different observables.

This review is devoted to the discussion of in-medium effects in the hadronic phase and its relation to the partial restoration of chiral symmetry. We, therefore, will not discuss another possible in-medium effect which is related to the deconfinement in the Quark Gluon Plasma, namely the suppression of the  $J/\Psi$ . This idea, which has been first proposed by Matsui and Satz [27] is based on the observation that due to the screening of the color interaction in the Quark Gluon Plasma, the  $J/\Psi$  is not bound anymore. As a result, if such a Quark Gluon Plasma is formed in a relativistic heavy ion collision, the abundance of  $J/\Psi$  should be considerably reduced. Of course also this signal suffers from more conventional backgrounds, namely the dissociation of the  $J/\Psi$  due to hadronic collisions [28, 29, 30]. To which extent present data can be understood in a purely hadronic scenario is extensively debated at the moment [31]. We refer the reader to the literature for further details [32].

In this review we will concentrate on proposed in-medium effects on hadrons. Whenever possible, we will try to disentangle conventional in-medium effects from those we believe are due to new physics, namely the change of the chiral condensate. We first will discuss medium effects on the Goldstone bosons, such as the pion, kaon and eta. Then we will concentrate on the vector mesons, which have the advantage that possible changes in their mass can be directly observed in the dilepton spectrum. We further discuss the effective mass of a nucleon in the nuclear medium. Finally, we will close by summarizing the current status of Lattice QCD calculations concerning hadronic properties. As will become clear from our discussion, progress in this field requires the input from experiment. Many question cannot be settled from theoretical consideration alone. Therefore, we will always try to emphasize the observational aspects for each proposed in-medium correction.

## 2 GOLDSTONE BOSONS

This first part is devoted to in-medium effects of Goldstone bosons, specifically the pion, kaon and eta. The fact that these particles are Goldstone bosons means that their properties are directly linked to the spontaneous breakdown of chiral symmetry. Therefore, rather reliable predictions about their properties can be made using chiral symmetry arguments and techniques, such as chiral perturbation theory. Contrary to naive expectations, the properties of Goldstone bosons are rather robust with respect to changes of the chiral condensate. On second

thought, however, this is not so surprising, because as long as chiral symmetry is spontaneously broken there will be Goldstone bosons. The actual nonvanishing value of their mass is due to the explicit symmetry breaking, which – at least in case of the pion – are small. Changes in their mass, therefore, are associated with the sub-leading explicit symmetry breaking terms and are thus small. We should stress, however, that in case of the kaon the symmetry breaking terms are considerably larger, leading to sizeable corrections to their mass at finite density as we shall discuss below.

## 2.1 *The pion*

Of all hadrons, the pion is probably the one where in-medium corrections are best known and understood. From the theoretical point of view, pion properties can be well determined because the pion is such a ‘good’ Goldstone boson. The explicit symmetry breaking terms in this case are small, as they are associated with the current masses of the light quarks. Therefore, the properties of the pion at finite density and temperature can be calculated using for instance chiral perturbation theory. But more importantly, there exists a considerable amount of experimental data ranging from pionic atoms to pion-nucleus experiments to charge-exchange reactions. These data all address the pion properties at finite density which we will discuss now.

### FINITE DENSITY

The mass of a pion in symmetric nuclear matter is directly related to the real part of the isoscalar-s-wave pion optical potential. This has been measured to a very high precision in pionic atoms [33], and one finds to first order in the nuclear density  $\rho$

$$\Delta m_\pi^2 = -4\pi(b_0)_{\text{eff}}\rho, \quad (b_0)_{\text{eff}} \simeq -0.024m_\pi^{-1}. \quad (1)$$

At nuclear matter density the shift of the pion mass is

$$\frac{\Delta m_\pi}{m_\pi} \sim 8\%, \quad (2)$$

which is small as expected from low energy theorems based on chiral symmetry. Theoretically, the value of  $(b_0)_{\text{eff}}$  is well understood as the combined contribution from a single scattering (impulse approximation) involving the small s-wave isoscalar  $\pi N$  scattering length  $b_0 =$

$-0.010(3)m_\pi^{-1}$  and the contribution from a (density dependent) correlation or re-scattering term, which is dominated by the comparatively large isovector-s-wave scattering length  $b_1 = -0.091(2)m_\pi^{-1}$ , i.e.,

$$(b_0)_{\text{eff}} = b_0 - [1 + \frac{m_\pi}{M_N}(b_0^2 + 2b_1^2)] < \frac{1}{r} >, \quad (3)$$

where the correlation length  $< 1/r > \simeq 3p_f/(2\pi)$  is essentially due to the Pauli exclusion principle. The weak and slightly repulsive s-wave pion potential has also been found in calculations based on chiral perturbation theory [34, 35] as well as the phenomenological Nambu - Jona-Lasinio model [36]. This is actually not too surprising because both amplitudes,  $b_0$  and  $b_1$ , are controlled by chiral symmetry and its explicit breaking [33, 37].

Pions at finite momenta, on the other hand, interact very strongly with nuclear matter through the p-wave interaction, which is dominated by the  $P_{33}$  delta-resonance. This leads to a strong mixing of the pion with nuclear excitations such as a particle-hole and a delta-hole excitation. As a result the pion-like excitation spectrum develops several branches, which to leading order correspond to the pion, particle-hole, and delta-hole excitations. Because of the attractive p-wave interaction the dispersion relation of the pionic branch becomes considerably softer than that of a free pion. A simple model, which has been used in practical calculations [38, 39] is the so-called delta-hole model [33, 40] which concentrates on the stronger pion-nucleon-delta interaction ignoring the nucleon-hole excitations. In this model, the pion dispersion relation in nuclear medium can be written as

$$\omega(\mathbf{k}, \rho) = m_\pi^2 + \mathbf{k}^2 + \Pi(\omega, \mathbf{k}), \quad (4)$$

where the pion self-energy is given by

$$\Pi(\omega, \mathbf{k}) = \frac{\mathbf{k}^2 \chi(\omega, \mathbf{k})}{1 - g' \chi(\omega, \mathbf{k})}, \quad (5)$$

with  $g' \approx 0.6$  the Migdal parameter which accounts for short-range correlations, also known as the Ericson-Ericson - Lorentz-Lorenz effect [41]. The pion susceptibility  $\chi$  is given by

$$\chi(\omega, \mathbf{k}) \approx \frac{8}{9} \left( \frac{f_{\pi N \Delta}}{m_\pi} \right)^2 \frac{\omega_R}{\omega^2 - \omega_R^2} \exp\left(-2\mathbf{k}^2/b^2\right) \rho, \quad (6)$$

where  $f_{\pi N \Delta} \approx 2$  is the pion-nucleon-delta coupling constant,  $b \approx 7m_\pi$  is the range of the form factor, and  $\omega_R \approx \frac{\mathbf{k}^2}{2m_\Delta} + m_\Delta - m_N$ .

The pion dispersion relation obtained in this model is shown in Fig. 1. The pion branch in the lower part of the figure is seen to become softened, while the delta-hole branch in the upper part of the figure is stiffened. Naturally this model oversimplifies things and once couplings to nucleon-hole excitations and corrections to the width of the delta are consistently taken into account, the strength of the delta-hole branch is significantly reduced [42, 43, 44, 45, 46, 47]. However, up to momenta of about  $2m_\pi$  the pionic branch remains pretty narrow and considerably softer than that of a free pion. The dispersion relation of the pion can be measured directly using  $(^3\text{He}, t)$  charge exchange reactions [48]. In these experiments the production of coherent pions have been inferred from angular correlations between the outgoing pion and the transferred momentum. The corresponding energy transfer is smaller than that of free pions indicating the attraction in the pion branch discussed above. More detailed measurements of this type are currently being analyzed [49].

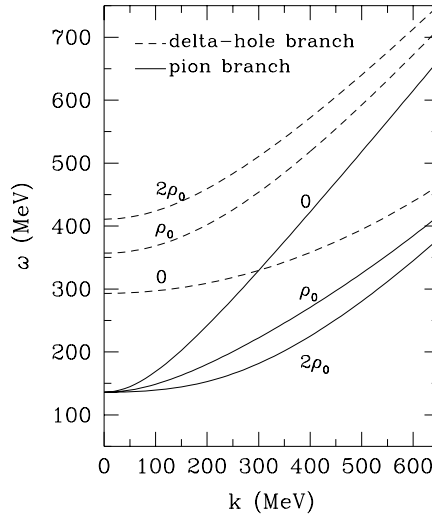


Figure 1: Pion dispersion relation in the nuclear medium. The normal nuclear matter density is denoted by  $\rho_0$ .

In the context of relativistic heavy ion collisions interest in the in-medium modified pion dispersion relation got sparked by the work of



Gale and Kapusta [50]. Since the density of states is proportional to the inverse of the group velocity of these collective pion modes, which can actually vanish depending on the strength of the interaction (see Fig. 1), they argued that a softened dispersion relation would lead to a strong enhancement of the dilepton yield. This happens at invariant mass close to twice the pion mass because the number of pions with low energy will be considerably increased. However, as pointed out by Korpa and Pratt [51] and further explored in more detail in [52], the initially proposed strong enhancement is reduced considerably once corrections due to gauge invariance are properly taken into account. Other places, where the effect of the pion dispersion relation is expected to play a role, are the inelastic nucleon-nucleon cross section [38] and the shape of the pion spectrum [39, 55, 56]. In the latter case one finds [39] that as a result of the attractive interaction the yield of pions at low  $p_t$  is enhanced by about a factor of 2. This low  $p_t$  enhancement is indeed seen in experimental data for  $\pi^-$  production from the BEVALAC [57] and in more recent data from the TAPS collaboration [58], which measures neutral pions. Certainly, in order to be conclusive, more refined calculations are needed, and they are presently being carried out [47].

#### FINITE TEMPERATURE

The pionic properties at finite temperature instead of finite density are qualitatively very similar, although the pion now interacts mostly with other pions instead of nucleons. Again, due to chiral symmetry, the s-wave interaction among pions is small and slightly repulsive, leading to a small mass shift of the pion [59].

And analogous to the case of nuclear matter, there is a strong attractive p-wave interaction, which is now dominated by the  $\rho$ -resonance. This again results in a softened dispersion relation which, however, is not as dramatic as that obtained at finite density [60, 61, 62, 63]. Also the phenomenological consequences are similar. The modified dispersion relation results in an enhancement of dileptons from the pion annihilation by a factor of about two at invariant masses around 300 – 400 MeV [64, 65, 66]. Unfortunately, in the same mass range other channels dominate the dilepton spectrum (see discussion in section 3.1) so that this enhancement cannot be easily observed in experiment [65]. The modified dispersion relation has also been invoked in order to explain the enhancement of low transverse momentum pions observed at CERN-SPS heavy ion collisions [61, 67, 68]. However, at these high energies the expansion velocity of the system, which is mostly made out of pions,

is too fast for the attractive pion interaction to affect the pion spectrum at low transverse momenta [62]. This is different at the lower energies around 1 GeV. There, the dominant part of the system and the source for the pion potential are the nucleons, which move considerably slower than pions. Therefore, pions have a chance to leave the potential well before it has disappeared as a result of the expansion.

To summarize this section on pions, the effects for the pions at finite density as well as finite temperature are dominated by p-wave resonances, the  $\Delta(1230)$  and the  $\rho(770)$ , respectively. Both are not related directly to chiral symmetry and its restoration but rather to what we call here many-body effects, which of course does not make them less interesting. The s-wave interaction is small because the pion is such a ‘good’ Goldstone Boson; probably too small to have any phenomenological consequence (at least for heavy ion collisions).

## 2.2 Kaons

Contrary to the pion, the kaon is not such a good Goldstone boson. Effects of the explicit chiral symmetry breaking are considerably bigger, as one can see from the mass of the kaon, which is already half of the typical hadronic mass scale of 1 GeV. In addition, since the kaon carries strangeness, its behavior in non-strange, isospin symmetric matter will be different from that of the pion. Rather interesting phenomenological consequences arise from this difference such as a possible condensation of antikaons in neutron star matter [69, 70, 71, 72].

### CHIRAL LAGRANGIAN

This difference can best be exemplified by studying the leading order effective  $SU(3)_L \times SU(3)_R$  Lagrangian obtained in heavy baryon chiral perturbation theory [72, 73]

$$\begin{aligned} \mathcal{L}_0 = & \frac{f^2}{4} \text{Tr} \partial^\mu U \partial_\mu U^\dagger + \frac{f^2}{2} r \text{Tr} M_q (U + U^\dagger - 2) \\ & + \text{Tr} \bar{B} i v_\mu \mathcal{D}^\mu B + 2D \text{Tr} \bar{B} S^\mu \{A_\mu, B\} \\ & + 2F \text{Tr} \bar{B} S^\mu [A_\mu, B]. \end{aligned} \quad (7)$$

In the above formula, we have  $U = \exp(2i\pi/f)$  with  $\pi$  and  $f$  being the pseudoscalar meson octet and their decay constant, respectively;  $B$  is the baryon octet;  $v_\mu$  is the four velocity of the heavy baryon ( $v^2 = 1$ ); and  $S_\mu$  stands for the spin operator  $S_\mu = \frac{1}{4}\gamma_5[v_\nu\gamma_\nu, \gamma_\mu]$  and  $v_\mu S^\mu = 0$ ;

$M_q$  is the quark mass matrix; and  $r$ ,  $D$  as well as  $F$  are empirically determined constants. Furthermore,

$$\mathcal{D}_\mu \mathcal{B} = \partial_\mu B + [V_\mu, B] \quad (8)$$

$$V_\mu = \frac{1}{2}(\xi \partial_\mu \xi^\dagger + \xi^\dagger \partial_\mu \xi), \quad A_\mu = \frac{i}{2}(\xi \partial_\mu \xi^\dagger - \xi^\dagger \partial_\mu \xi), \quad (9)$$

with  $\xi^2 = U$ .

To leading order in the chiral counting the explicit symmetry breaking term  $\sim \text{Tr } M_q(U + U^\dagger - 2)$  gives rise to the masses of the Goldstone bosons. The interesting difference between the behavior of pions and kaons in matter arises from the term involving the vector current  $V_\mu$ , i.e.  $\text{Tr } \bar{B} i v_\mu \mathcal{D}^\mu B$ . In case of the pion, this term is identical to the well-known Weinberg-Tomozawa term [37]

$$\delta \mathcal{L}_{WT} = \frac{-1}{4f_\pi^2} (\bar{N} \vec{\tau} \gamma^\mu N) \cdot (\vec{\pi} \times (\partial_\mu \vec{\pi})). \quad (10)$$

It contributes only to the isovector s-wave scattering amplitude and, therefore, does not contribute to the pion optical potential in isospin symmetric nuclear matter. This is *different* in case of the kaon, where we have

$$\delta \mathcal{L}_{WT} = \frac{-i}{8f^2} \left( 3(\bar{N} \gamma^\mu N)(K \overset{\leftrightarrow}{\partial}_\mu K) + (\bar{N} \vec{\tau} \gamma^\mu N)(K \vec{\tau} \overset{\leftrightarrow}{\partial}_\mu K) \right). \quad (11)$$

The first term contributes to the isoscalar s-wave amplitude and, therefore, gives rise to an attractive or repulsive optical potential for  $K^-$  and  $K^+$  in symmetric nuclear matter. It turns out, however, that this leading order Lagrangian leads to an s-wave scattering length which is too repulsive as compared with experiment. Therefore, terms next to leading order in the chiral expansion are needed. Some of these involve the kaon-nucleon sigma term and thus are sensitive to the second difference between pions and kaons, namely the strength of the explicit symmetry breaking. The next to leading order effective kaon-nucleon Lagrangian can be written as [72]

$$\begin{aligned} \mathcal{L}_{\nu=2} = & \frac{\Sigma_{KN}}{f^2} (\bar{N} N) (\bar{K} K) + \frac{C}{f^2} (\bar{N} \vec{\tau} N) \cdot (\bar{K} \vec{\tau} K) \\ & + \frac{\tilde{D}}{f^2} (\bar{N} N) (\partial_t \bar{K} \partial_t K) + \frac{\tilde{D}'}{f^2} (\bar{N} \vec{\tau} N) \cdot (\partial_t \bar{K} \vec{\tau} \partial_t K). \end{aligned} \quad (12)$$

The value of the kaon-nucleon sigma-term  $\Sigma_{KN} = \frac{1}{2}(m_q + m_s) \langle N | \bar{u}u + \bar{s}s | N \rangle$  depends on the strangeness content of the nucleon,  $y = 2 \langle N | \bar{s}s | N \rangle$

$\langle N|\bar{u}u + \bar{d}d|N\rangle \approx 0.1 - 0.2$ . Using the light quark mass ratio  $m_s/m \approx 29$ , one obtains  $370 < \Sigma_{KN} < 405$  MeV. The additional parameters,  $C, \tilde{D}, \tilde{D}'$  are then fixed by comparing with  $K^+$ -nucleon scattering data [72]. This so determined effective Lagrangian can then be used to predict the  $K^-$ -nucleon scattering amplitudes, and one obtains an attractive isoscalar s-wave scattering length in contradiction with experiments, where one finds a repulsive amplitude [74]. This discrepancy has been attributed to the existence of the  $\Lambda(1405)$  which is located below the  $K^-N$ -threshold. From the analysis of  $K^-p \rightarrow \Sigma\pi$  reactions it is known that this resonance couples strongly to the  $I = 0$   $K^-p$  state, and, therefore, leads to repulsion in the  $K^-p$  amplitude.

#### THE ROLE OF THE $\Lambda(1405)$

Already in the sixties [75] there have been attempts to understand the  $\Lambda(1405)$  as a bound state of the proton and the  $K^-$ . In this picture, the underlying  $K^-$ -proton interaction is indeed attractive as predicted by the chiral Lagrangians but the scattering amplitude is repulsive only because a bound state is formed. This concept is familiar to the nuclear physicist from the deuteron, which is bound, because of the attractive interaction between proton and neutron. The existence of the deuteron then leads to a repulsive scattering length in spite of the attractive interaction between neutron and proton. Chiral perturbation theory is based on a systematic expansion of the S-matrix elements in powers of momenta and, therefore, effects which are due to the proximity of a resonance, such as the  $\Lambda(1405)$ , will only show up in terms of rather high order in the chiral counting. Thus, it is not too surprising that the first two orders of the chiral expansion predict the wrong sign of the  $K^-N$  amplitude. To circumvent this problem, the chiral perturbation calculation has been extended to either include an explicit  $\Lambda(1405)$  state [76] or to use the interaction obtained from the leading order chiral Lagrangian as a kernel for a Lippman-Schwinger type calculation, which is then solved to generate a bound state  $\Lambda(1405)$  [77].

This picture of the  $\Lambda(1405)$  as a  $K^-p$  bound state has recently received some considerable interest in the context of in-medium corrections. In ref. [78] it has been pointed out that in this picture as a result of the Pauli-blocking of the proton inside this bound state, the properties of the  $\Lambda(1405)$  would be significantly changed in the nuclear environment. With increasing density, its mass increases and the strength of the resonance is reduced (see Fig. 2). Because of this shifting and ‘disappearance’ of the  $\Lambda(1405)$  in matter, the  $K^-$  optical potential changes

sign from repulsive to attractive at a density of about 1/4 of nuclear matter density (see Fig. 2) in agreement with a recent analysis of  $K^-$  atoms [79]. These findings have been confirmed in ref. [80]. This in-medium change of the  $\Lambda(1405)$  due to the Pauli blocking can only occur if a large fraction of its wave function is indeed that of a  $K^-$ -proton bound state. Of course one could probably allow for a small admixture of a genuine three quark state without changing the results for the measured kaon potentials. But certainly, if the  $\Lambda(1405)$  is mostly a genuine three quark state, the Pauli blocking should not affect its properties. Therefore, it would be very interesting to confirm the mass shift of the  $\Lambda(1405)$  for instance by a measurement of the missing mass spectrum of kaons in the reaction  $p + \gamma \rightarrow \Lambda(1405) + K^+$  at CEBAF. Thus, the atomic nucleus provides a unique laboratory to investigate the properties of elementary particles. From our discussion it is clear that this in-medium change of the  $\Lambda(1405)$  is not related to the restoration of chiral symmetry.

#### EXPERIMENTAL RESULTS

Phenomenologically, the attractive optical potential for the  $K^-$  in nuclear matter is of particular interest because it can lead to a possible kaon condensation in neutron stars [72, 81]. This would limit the maximum mass of neutron stars to about one and a half solar masses and give rise to speculations about many small black holes in our galaxy [82]. Kaonic atoms, of course, only probe the very low density behavior of the kaon optical potential and, therefore, an extrapolation to the large densities relevant for neutron stars is rather uncertain. Additional information about the kaon optical potential can be obtained from heavy ion collision experiments, where densities of more than twice nuclear matter density are reached.

Observables that are sensitive to the kaon mean-field potentials are the subthreshold production [83, 84, 85] as well as kaon flow [86, 87]. Given an attractive/repulsive mean-field potential for the kaons it is clear that the subthreshold production is enhanced/reduced. In case of the kaon flow an attractive interaction between kaons and nucleons aligns the kaon flow with that of the nucleons whereas a repulsion leads to an anti-alignment (anti-flow). However, both observables are also extremely sensitive to the overall reaction dynamics, in particular to the properties of the nuclear mean field and to reabsorption processes especially in case of the antikaons. Therefore, transport calculations are required in order to consistently incorporate all these effects.

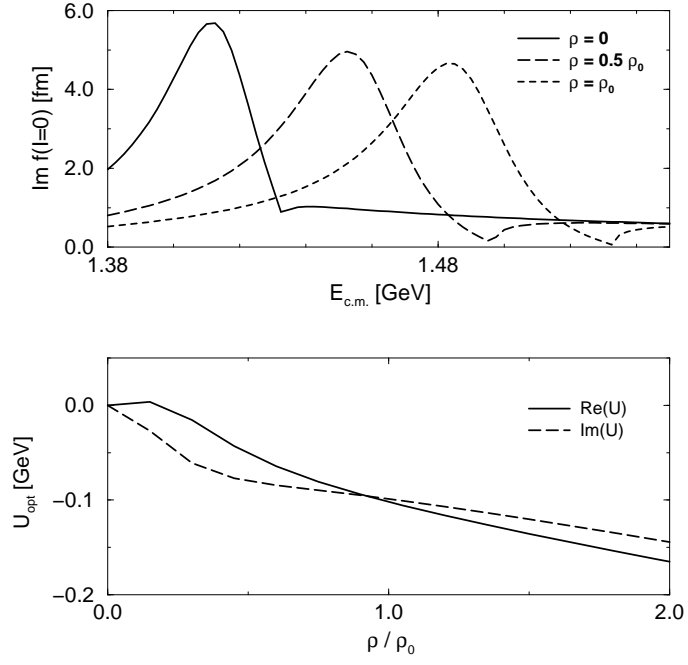


Figure 2: (a) Imaginary part of the  $I = 0$   $K^-$ -proton scattering amplitude for different densities. (b) real and Imaginary part of the  $K^-$  optical potential.

In Fig. 3 we show the result obtained from such calculations for the  $K^+$  subthreshold production and flow together with experimental data (solid circles) from the KaoS collaboration [88] and from the FOPI collaboration [89] at GSI. Results are shown for calculations without any mean fields (dotted curves) as well as with a repulsive mean field obtained from the chiral Lagrangian (solid curves), which is consistent with a simple impulse approximation. The experimental data are nicely reproduced when the kaon mean-field potential is included. Although the subthreshold production of kaons can also be explained without any kaon mean field [90, 91], the assumption used in these calculations that a lambda particle has the same mean-field potential as a nucleon is not consistent with the phenomenology of hypernuclei [92]. Since it is undisputed that these observables are sensitive to the in-medium kaon potential, a systematic investigation including all observables should eventually reveal more accurately the strength of the kaon potential in dense matter. Additional evidence for a repulsive  $K^+$  potential comes from  $K^+$ -nucleus experiments. There the measured cross sections and angular distributions can be pretty well understood within a simple impulse approximation. Actually it seems that an additional (15%) repulsion is required to obtain an optimal fit to the data [93], which has been suggested as a possible evidence for a swelling of the nucleon size or a lowering of the omega meson mass in the nuclear medium [94].

As for the  $K^-$ , both chiral perturbation theory and dynamical models of the  $K^-$ -nucleon interaction [78] indicate the existence of an attractive mean field, so the same observables can be used [95, 96]. However, the measurement and interpretation of  $K^-$  observables is more difficult since it is produced less abundantly. Also reabsorption effects due to the reaction  $K^- N \rightarrow \Lambda \pi$  are strong, which further complicates the analysis. Nevertheless, the effect of the attractive mean field has been shown to be significant as illustrated in Fig. 4, where results from transport calculations with (solid curves) or without (dotted curves) attractive mean field for the antikaons [95] are shown. It is seen that the data (solid circles) on subthreshold  $K^-$  production [97] support the existence of an attractive antikaon mean-field potential. For the  $K^-$  flow, there only exist very preliminary data from the FOPI collaboration [98], which seem to show that the  $K^-$ 's have a positive flow rather an antiflow, thus again consistent with an attractive antikaon mean-field potential.

Let us conclude this section on kaons by pointing out that the properties of kaons in matter are qualitatively described in chiral perturbation theory. Although some of the effects comes from the Weinberg-Tomozawa vector type interaction, which contributes because contrary

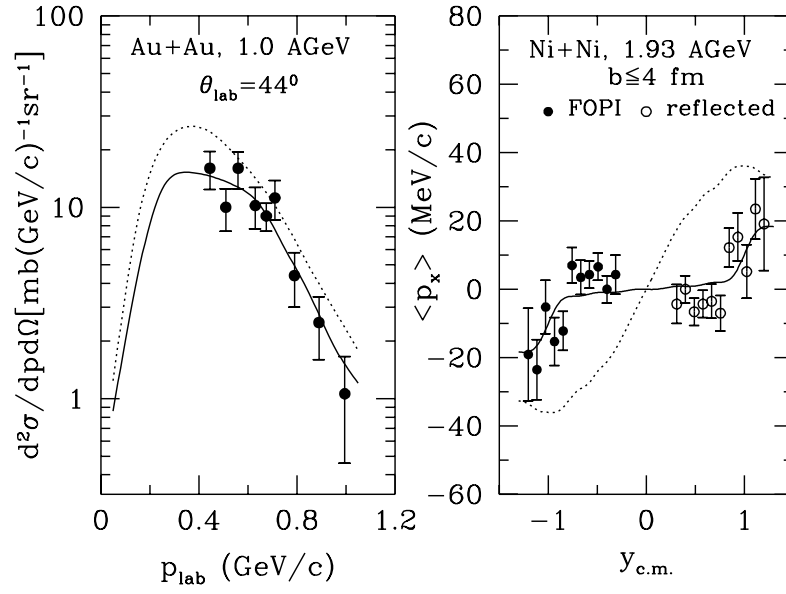


Figure 3: Kaon yield (left panel) and flow (right panel) in heavy ion collisions. The solid and dotted curves are results from transport model calculations with and without kaon mean-field potential, respectively. The data are from refs. [88, 89].



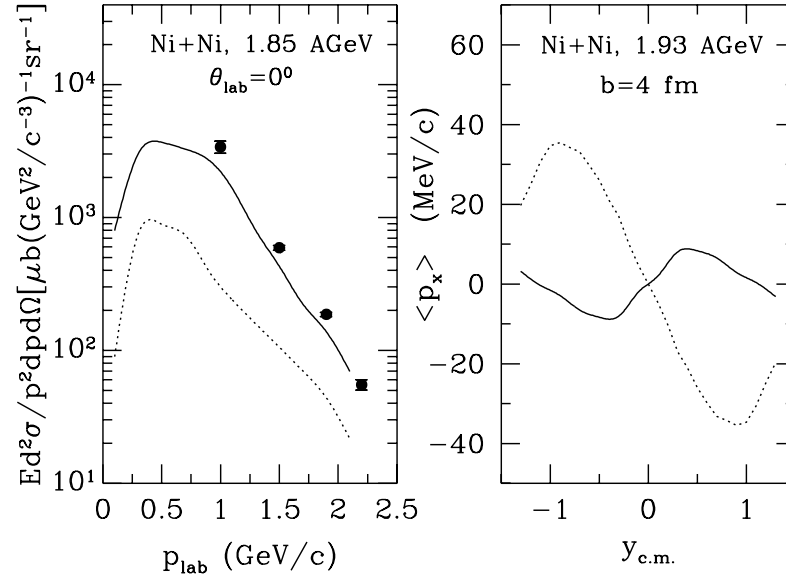


Figure 4: Antikaon yield (left panel) and flow (right panel) in heavy ion collisions. The solid and dotted curves are results from transport model calculations with and without kaon mean-field potential, respectively. The data are from ref. [97].

to the pion the kaon has only one light quark, in order to explain the experimental data on kaon subthreshold production and flow in heavy ion collisions an additional attractive scalar interaction is required, which is related to the explicit breaking of chiral symmetry and higher order corrections in the chiral expansion.

### 2.3 *Etas*

The properties of etas are not only determined by consideration of chiral symmetry but also by the explicit breaking of the  $U_A(1)$  axial symmetry due to the axial anomaly in QCD. As a result, the singlet eta, which would be a Goldstone boson if  $U_A(1)$  were not explicitly broken, becomes heavy. Furthermore, because of  $SU(3)$  symmetry breaking – the strange quark mass is considerably heavier than that of up and down quark – the octet eta and singlet eta mix, leading to the observed particles  $\eta$  and  $\eta'$ . The mixing is such that the  $\eta'$  is mostly singlet and thus heavy, and the  $\eta$  is mostly octet and therefore has roughly the mass of the kaons. If, as has been speculated [99, 100], the  $U_A(1)$  symmetry is restored at high temperature due to the instanton effects, one would expect considerable reduction in the masses of the etas as well as in their mixing [101, 102]. A dropping eta meson in-medium mass is expected to provide a possible explanation for the observed enhancement of low transverse momentum etas in SIS heavy ion experiments at subthreshold energies [103]. However, an analysis of photon spectra from heavy ion collisions at SPS-energies puts an upper limit on the  $\eta/\pi^0$  ratio to be not more than 20 % larger for central than for peripheral collisions [104]. This imposes severe constraints on the changes of the  $\eta$  properties in hadronic matter.

We note that the restoration of chiral symmetry is important in the  $\eta - \eta'$  sector, because without the  $U_A(1)$  breaking both would be Goldstone bosons of an extended  $U(3) \times U(3)$  symmetry. However, it is still being debated what the effects precisely are.

## 3 VECTOR MESONS

Of all particles it is probably the  $\rho$ -meson which has received the most attention in regards of in-medium corrections. This is mainly due to the fact that the  $\rho$  is directly observable in the dilepton invariant mass spectrum. Also, since the  $\rho$  carries the quantum numbers of the conserved vector current, its properties are related to chiral symmetry and can,

as we shall discuss, be investigated using effective chiral models as well as current algebra and QCD sum rules. That possible changes of the  $\rho$  can be observed in the dilepton spectrum measured in heavy ion collisions has been first demonstrated in ref. [105] and then studied in more detail in [106, 107]. The in-medium properties of the  $a_1$  are closely related to that of the  $\rho$  since they are chiral partners and their mass difference in vacuum is due to the spontaneous breaking of chiral symmetry [108]. Unfortunately, there is no direct method to measure the changes of the  $a_1$  in hadronic matter. The  $\omega$ -meson on the other hand is a chiral singlet, and the relation of its properties to chiral symmetry is thus not so direct. However, calculations based on QCD sum rules predict also changes in the  $\omega$  mass as one approaches chiral restoration. Observationally, the  $\omega$  can probably be studied best due to its rather small width and its decay channel into dileptons. Finally, there is the  $\phi$  meson. In an extended SU(3) chiral symmetry, the  $\phi$  and the  $\omega$  are a superposition of the singlet and octet states with nearly perfect mixing, i.e. the  $\phi$  contains only strange quarks whereas the  $\omega$  is made only out of light quarks. Both QCD sum rules [109] and effective chiral models [110] predict a lowering of the  $\phi$ -meson mass in medium.

The question on whether or not masses of light hadrons change in the medium has received considerable interest as a result of the conjecture of Brown and Rho [8], which asserts that the masses of all mesons, with the exception of the Goldstone Bosons, should scale with the quark condensate. While the detailed theoretical foundations of this conjecture are still being worked on [111, 112, 113] the basic argument of Brown and Rho is as follows. Hadron masses, such as that of the  $\rho$ -meson, violate scale invariance, which is a symmetry of the classical QCD Lagrangian. In QCD scale invariance is broken on the quantum level by the so-called trace anomaly (see e.g. [114]), which is proportional to the Gluon condensate. Thus one could imagine that with the disappearance of the gluon condensate, i.e. the bag pressure, scale invariance is restored, which on the hadronic level implies that hadron masses have to vanish. Therefore, one could argue that hadron masses should scale with the bag-pressure as originally proposed by Pisarski [115]. But the conjecture of Brown and Rho goes even further. They assume that the gluon condensate can be separated into a hard and soft part, the latter of which scales with the quark condensate and is also responsible for the masses of the light hadrons. This picture finds some support from lattice QCD calculations in that the gluons condensate drops by about 50% at the chiral phase transition [116]. To what extent this is also reflected in changes in hadron masses is not clear at the moment

(see section 5), although one should mention that the rise in the entropy density close to the critical temperature can be explained if one assumes the hadron masses to scale with the quark condensate [116]. Another aspect of the Brown and Rho scaling is that once the scaling hadron masses are introduced in the chiral Lagrangian, only tree-level diagrams are needed as the contribution from higher order diagrams is expected to be suppressed. In their picture, Goldstone bosons are not subject to this scaling, since they receive their mass from the explicit chiral symmetry breaking due to finite current quark masses, which are presumably generated at a much higher (Higgs) scale.

### 3.1 *The rho meson*

Since the  $\rho$  is a vector meson, it couples directly to the isovector current which then results in the direct decay of the  $\rho$  into virtual photons, i.e. dileptons. Consequently, properties of the  $\rho$  meson can be investigated by studying two-point correlation functions of the isovector currents, i.e.,

$$\Pi_{\mu\nu}(q) = i \int e^{iqx} \langle T J_\mu(x) J_\nu(0) \rangle_\rho d^4x. \quad (13)$$

The masses of the rho meson and its excitations ( $\rho' \dots$ ) correspond to the positions of the poles of this correlation function. This is best seen if one assumes that the current field identity [117] holds, namely that the current operator is proportional to the  $\rho$ -meson field. In this case the above correlation function is identical, up to a constant, to the  $\rho$ -meson propagator. The imaginary part of this correlation function is also directly proportional to the electron-positron annihilation cross section [118], where the  $\rho$ -meson is nicely seen. In addition, at higher center-of-mass energies, one sees a continuum in the electron-positron annihilation cross section, which corresponds to the excited states of the rho mesons as well as to the onset of perturbative QCD processes.

In-medium changes of the  $\rho$  meson can be addressed theoretically by evaluating the current-current correlator in the hadronic environment. The current-current correlator can be evaluated either in effective chiral models, or using current algebra arguments, or directly in QCD. In the latter, one evaluates the correlator in the deeply Euclidean region ( $q^2 \rightarrow -\infty$ ) using the Wilson expansion, where all the long distance physics is expressed in terms of vacuum expectation values of quark and gluon operators, the so-called condensates. Dispersion relations are then

used to relate the correlator in the Euclidean region to that for positive  $q^2$ , where the hadronic eigenstates are located. One then assumes a certain shape for the phenomenological spectral functions, typically a delta function, which represents the bound state, and a continuum, which represents the perturbative regime. These so called QCD sum rules, therefore, relate the observable hadronic spectrum with the QCD vacuum condensates (for a review of the QCD sum-rule techniques see e.g. [119]). These relations can then either be used to determine the condensates from measured hadronic spectra, or, to make predictions about changes of the hadronic spectrum due to in-medium changes of the condensates.

Similarly, one can study the properties of  $a_1$  meson in the nuclear medium through the axial vector correlation function. Then, once chiral symmetry is restored, there should be no observable difference between left-handed and right-handed or equivalently vector and axial vector currents. Consequently, the vector and axial vector correlators should be identical. Often, this identity of the correlators is identified with the degeneracy of the  $\rho$  and  $a_1$  mesons in a chirally symmetric world. This, however, is not the only possibility, as was pointed out by Kapusta and Shuryak [120]. There are at least three qualitatively different scenarios, for which the vector and axial vector correlator are identical (see Fig. 5).

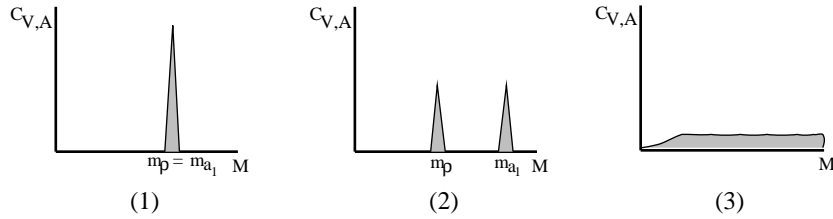


Figure 5: Several possibilities for the vector and axial-vector spectral functions in the chirally restored phase.

1. The masses of  $\rho$  and  $a_1$  are the same. The value of the common mass, however, does not follow from chiral symmetry arguments alone.
2. The mixing of the spectral functions, i.e, both the vector and axial-vector spectral functions have peaks of similar strength at both the mass of the  $\rho$  and the mass of the  $a_1$ .

3. Both spectral functions could be smeared over the entire mass range. Because of thermal broadening of the mesons and the onset of deconfinement, the structure of the spectral function may be washed out, and it becomes meaningless to talk about mesonic states.

#### FINITE TEMPERATURE

At low temperatures, where the heat bath consists of pions only, one can employ current algebra as well as PCAC to obtain an essentially model-independent result for the properties of the  $\rho$ . Using this technique Dey, Eletzky and Ioffe [121] could show that to leading order in the temperature,  $T^2$ , the mass of the  $\rho$ -meson does not change. Instead one finds an admixture of the axial-vector correlator, i.e. that governed by the  $a_1$ -meson. Specifically to this order the vector correlator is given by

$$C_V(q, T) = (1 - \epsilon) C_V(q, T = 0) + \epsilon C_A(q, T = 0) + \mathcal{O}(T^4), \quad (14)$$

with  $\epsilon = T^2/(6f_\pi^2)$ . Here  $C_V, C_A$  stand for the vector-isovector and axial-vector-isovector correlation functions, respectively. One should note, that to the same order the chiral condensate is reduced [59],

$$\frac{\langle \bar{q}q \rangle_T}{\langle \bar{q}q \rangle_0} = 1 - \frac{T^2}{8f_\pi^2} + \mathcal{O}(T^4). \quad (15)$$

Therefore, to leading order a drop in the chiral condensate does not affect the mass of the  $\rho$ -meson but rather reduces its coupling to the vector current and in particular induces an admixture of the  $a_1$ -meson. This finding is at variance with the Brown-Rho scaling hypothesis. The reason for this difference is not yet understood. The admixture of the axial correlator is directly related with the onset of chiral restoration. If chiral symmetry is restored, the vector and axial-vector correlators should be identical. The result of Dey et al. suggests that this is achieved by a mixing of the two instead of a degeneracy of the  $\rho$  and  $a_1$  masses. If the mixing is complete, i.e.  $\epsilon = 1/2$ , then the extrapolation of the low temperature result (14) would give a critical temperature of  $T_c = \sqrt{3}f_\pi \simeq 164 \text{ MeV}$ , which is surprisingly close to the value given by recent lattice calculations.

Corrections to the order  $T^4$  involve physics beyond chiral symmetry. As nicely discussed in ref. [122], to this order the contributions can be separated into two distinct contributions. The first arises if a pion from the heat bath couples via a derivative coupling to the current

under consideration. In this case the contribution is proportional to the invariant pion density

$$n_\pi = \int \frac{d^3q}{2\omega(2\pi)^3} e^{-\omega/T} \sim T^2, \quad (16)$$

and the square of the pion momentum  $q^2 \sim T^2$ . A typical example would be for instance the self-energy correction to the  $\rho$  meson from the standard two-pion loop diagram involving the p-wave  $\pi\pi\rho$  coupling. The other contribution comes from interactions of two pions from the heat bath with the current, without derivative couplings. This is proportional to the square of the invariant pion density and thus  $\sim T^4$  as well. These pure density contributions again can be evaluated in a model independent fashion using current algebra techniques and, as before, do not change the mass of the  $\rho$ -meson, but change the coupling to the current and induce a mixing with the axial-vector correlator. Actually, it can be shown [122] that to all orders in the pion density these pure density effects do not change the  $\rho$  mass but induce mixing and reduce the coupling. At the same time, however, these contributions reduce the chiral condensate.

The contributions due to the finite pion momenta have been estimated in [122] to give a downward shift of both the  $\rho$  and  $a_1$  mass of about

$$\frac{\delta m}{m} \simeq 10\%. \quad (17)$$

for temperatures of  $\sim 150 - 200$  MeV. We should point out, however, that in this analysis those changes in the masses arise from Lorentz-nonscalar condensates in the operator product expansion. Therefore, they are not directly related to the change of the Lorentz-scalar chiral-condensate.

Effective chiral models also have been used to explore the order  $T^4$  corrections. Ref. [123] reports that the mass of the  $\rho$  drops whereas that of the  $a_1$  increases to this order. This is somewhat at variance with the findings of [124] where no drop in the mass of the  $\rho$  but a decrease in the  $a_1$  mass has been found. In this calculation, however, the terms leading to the order  $T^4$  changes have not been explicitly identified but rather a calculation to one loop order has been carried out. Presently, the difference in these results is not understood. Also the difference between the analysis of [122] and the effective chiral models is not resolved yet. The analysis of [122] uses dispersion relations to relate phenomenological space-like photon-pion amplitudes with the time-like ones needed to

calculate the mass shift. The effective Lagrangian methods, on the other hand, rely on pion-pion scattering data as well as measured decay width in order to fix their model parameters. One would think that both methods should give a reasonably handle on the leading order momentum-dependent couplings.

At higher temperatures as well as at finite density vacuum effects due to the virtual pair correction or the nucleon-antinucleon polarization could become important and they tend to reduce vector meson masses [125]. Because of the large mass of the nucleon, these corrections however, do not affect the leading temperature result  $\sim T^2$  discussed previously. Also this approach assumes that the physical vacuum consists of nucleon-antinucleon rather than quark-antiquark fluctuations. Whether this is the correct picture is, however, not yet resolved.

#### FINITE DENSITY

Since the density effect on the chiral condensate is much stronger than that of the temperature, one expects the same for the rho meson mass. But the situation is more complicated at finite density as one cannot make use of current algebra arguments and thus model-independent result as the one discussed previously are not available at this time. Present model calculations, however, disagree even on the sign of the mass shift. One class of models [126, 127, 128, 129, 130] considers the  $\rho$  as a pion-pion resonance and calculates its in-medium modifications due to those for the pions as discussed in section 2.1. These calculations typically show an increased width of the rho since it now can also decay into pion-nucleon-hole or pion-delta-hole states. At the same time this leads to an increased strength below the rho meson mass. The mass of the rho meson, defined as the position where the real part of the correlation function goes through zero, is shifted only very little. Most calculations give an upward shift but also a small downward shift has been reported [130, 131]. This difference seems to depend on the specific choice of the cutoff functions for the vertices involved [132], and thus is model dependent. However, the imaginary part of the correlation function, the relevant quantity for the dilepton measurements, hardly depends on a small upward or downward shift of the rho. The important and apparently model independent feature is the increased strength at low invariant masses due to the additional decay-channels available in nuclear matter.

Calculation using QCD-sum rules [133] predict a rather strong decrease of the  $\rho$ -mass with density ( $\sim 20\%$  at nuclear matter density),



which is similar to the much discussed Brown-Rho scaling [8]. Here, as in the finite temperature case, the driving term is the four quark condensate which is assumed to factorize

$$\langle (\bar{q}\gamma_\mu\lambda^a q)(\bar{q}\gamma^\mu\lambda^a q) \rangle_\rho \approx -\frac{16}{9}\langle\bar{q}q\rangle_\rho^2. \quad (18)$$

To which extent this factorization is correct at finite density is not clear. Also, when it comes to the parameterization of the phenomenological spectral distribution, these calculations usually assume the standard pole plus continuum form with the addition of a so called Landau damping contribution at  $q^2 = 0$  [133]. The additional strength below the mass of the rho as predicted by the previously discussed models is usually ignored. An attempt, however, has been made to see if a spectral distribution obtained from the effective models described above does saturate the QCD sum rules [134]. These authors found that they could only saturate the sum rule if they assumed an additional mass shift of the rho-meson peak downwards by  $\sim 140$  MeV. However, in this calculation only leading order density corrections to the quark condensates whereas infinity order density effects have been taken into account in order to calculate the spectrum. Another comparison with the QCD sum rules has recently been carried out in ref. [131] where a reasonable saturation of the sum-rule is reported using a similar model for the in medium correlation function.

There are also attempts to use photon-nucleon data in order to estimate possible mass shifts of the  $\rho$  and  $\omega$  mesons. In ref. [135] a simple pion and sigma exchange model is used in order to fit data for photoproduction of  $\rho$ - and  $\omega$ -mesons. Assuming vector dominance, this model is then used to calculate the self-energy of these vector mesons in nuclear matter. The authors find a downward shift of the  $\rho$  of about 18% at nuclear matter density, in rough agreement with the prediction from QCD sum rules. Quite to the contrary, ref. [136] using photoabsorption data and dispersion relations find an upwards shift of 10 MeV or 50 MeV for the longitudinal and transverse part of the  $\rho$ , respectively. This result, however, is derived for a  $\rho$ -meson which is not at rest in the nuclear matter frame, and, therefore a direct comparison of the two predictions is not possible.

Very recently Friman and Pirner have pointed out that the  $\rho$ -meson couples very strongly with the  $N^*(1720)$  resonance [31, 137]. The coupling is of p-wave nature and, similarly to the pion coupling to delta-hole states, the  $\rho$  may mix with  $N^*$ -nucleon-hole states resulting in a modified dispersion relation for the  $\rho$ -meson in nuclear matter. Since

the coupling is p-wave, only  $\rho$ -mesons with finite momentum are modified and shifted to lower masses. This momentum dependence of the low mass enhancement in the dilepton spectrum is an unique prediction which can be tested in experiment.

#### EXPERIMENTAL RESULTS

First measurement of dileptons in heavy ion collisions have been carried out by the DLS collaboration [53, 54] at the BEVALAC. The first published data based on a limited data set could be well reproduced using state of the art transport models [106, 138, 139] without any additional in-medium corrections. However, a recent reanalysis [140] including the full data set seems to show a considerable increase over the originally published data. It remains to be seen if these new data can also be understood without any in-medium corrections to hadronic properties. Recently, low mass dilepton data taken at the CERN SPS have been

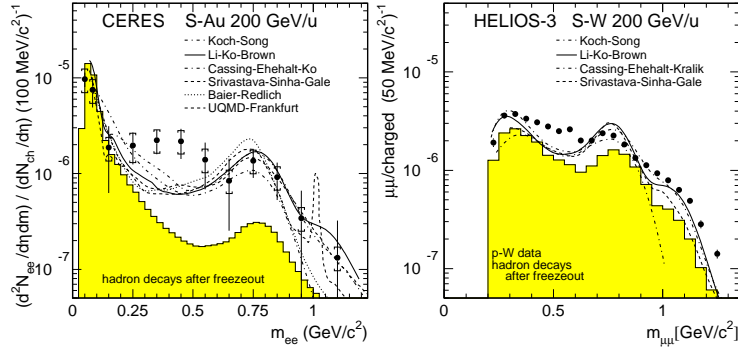


Figure 6: Dilepton invariant mass spectrum from S+Au and S+W collisions at 200 GeV/nucleon without dropping vector meson masses. The experimental statistical errors are shown as bars and the systematic errors are marked independently by brackets. (The figure is from ref. [141])

published by the CERES collaboration [142]. Their measurements for p+Be and p+Au could be well understood within a hadronic cocktail, which takes into account the measured particle yields from p+p experiments and their decay channels into dileptons. In case of the heavy

ion collision, S+Pb, however, the hadronic cocktail considerably under-predicted the measured data in particular in the invariant mass region of  $300 \text{ MeV} \leq M_{\text{inv}} \leq 500 \text{ MeV}$ . A similar enhancement has also been reported by the HELIOS-3 collaboration [143].

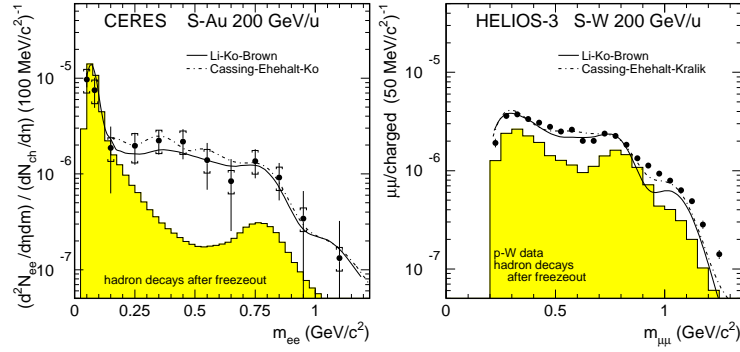


Figure 7: Same as Fig. 6 with dropping vector meson masses. (The figure is from ref. [141])

Of course it is well-known that in a heavy ion collisions at SPS-energies a hadronic fireball consisting predominantly of pions is created. These pions can pairwise annihilate into dileptons giving rise to an additional source, which has not been included into the hadronic cocktail. While the pion annihilation contributes to the desired mass range, many detailed calculations [65, 144, 145], which have taken this channel into account, still underestimate the data by about a factor of three as shown Fig. 6. Also, additional hadronic processes, such as  $\pi\rho \rightarrow \pi e^+e^-$ , have been taken into account [146, 147]. However, ‘conservative’ calculations could at best reach the lower end of the sum of statistical and systematic error bars of the CERES data. Furthermore, in-medium corrections to the pion annihilation process together with corrections to the  $\rho$ -meson spectral distribution have been considered using effective chiral models [64, 148]. But, these in-medium corrections, while enhancing the contribution of the pion annihilation channel somewhat, are too small to reproduce the central data points. Only models which allow for a dropping of the  $\rho$ -meson mass give enough yield in the low mass region [144, 149, 150, 151] as shown in Fig. 7. In these models, the change

of the rho meson mass is obtained from either the Brown-Rho scaling [149, 150, 151] or the QCD sum rules [144]. Furthermore, the vector dominance is assumed to be suppressed in the sense that the pion electromagnetic form factor, which in free space is dominated by the rho meson and is proportional to the square of the rho meson mass, is reduced as a result of the dropping rho meson mass. At finite temperature, such a suppression has been shown to exist in both the hidden gauge theory [152] and the perturbative QCD [153], where the pion electromagnetic form factor is found to decrease at the order  $T^2$ . This effect is related to the mixing between the vector-isovector and axial-vector-isovector correlators at finite temperature we discussed earlier. In most studies of dilepton production without a dropping rho meson mass, this effect has not been included. Although the temperature reached in heavy ion collisions at SPS energies is high and the ratio of pions to baryons in the final state is about 5 to 1, the authors in [144, 149, 150, 151] found that they had to rely on the nuclear density rather than on the temperature in order to reproduce the data.

There are also preliminary data from the Pb+Au collisions at 150 GeV/nucleon [141], which also seem to be consistent with the dropping in-medium rho meson scenario as well. Unfortunately, the experimental errors in this case are even larger than in the S+Au data and thus these new, preliminary data do not further discriminate between the dropping rho-mass and more conventional scenarios. Thus, additional data with reduced error bars are needed before firm conclusions about in-medium changes of the rho meson mass can be drawn. Although the data from the HELIOS-3 collaboration shown in the right panels of Figs. 6 and 7 do have very small errors, these are only the statistical ones while the systematic errors are not known.

To summarize this section on the rho meson, our present theoretical understanding for possible in-medium changes at low temperatures and vanishing baryon density indicates that one expects no (to leading order) and a small (to next to leading order) changes of the  $\rho$ -mass. At the same time to leading order in the temperature, the chiral condensate is reduced. Therefore, a direct connection between changes in the chiral condensate and the mass of the  $\rho$  meson has not been established as assumed in the Brown-Rho scaling. However, it may well be that the latter is only valid at higher temperature near the chiral phase transition. At finite density additional effects such as nuclear many-body excitations come into play, which make the situation much more complicated. Nevertheless, the observed enhancement of low mass dileptons from CERN-SPS heavy ion collisions can best be described by a

dropping rho meson mass in dense matter. Finite density and zero temperature is another area where possible changes of the vector mesons can be measured in experiment. Photon, proton and pion induced dilepton production from nuclei will soon be measured at CEBAF [154] and at GSI [155]. In principle these measurements should be able to determine the entire spectral distribution. If the predictions of the Brown-Rho scaling and the QCD sum rules are correct, mass shifts of the order of 100 MeV should occur, which would be visible in these experiments. Furthermore, by choosing appropriate kinematics, the properties of the  $\rho$  meson at rest as well as at finite momentum with respect to the nuclear rest frame can be measured.

### 3.2 *The omega meson*

The omega meson couples to the isoscalar part of the electromagnetic current. In QCD sum rules it is also dominated by the four quark condensate. If the latter is reduced in medium, then the omega meson mass also decreases. Assuming factorization for the four quark condensate, Hatsuda and Lee [133] showed that the change of the omega meson mass in dense nuclear matter would be similar to that of the rho meson mass. The dropping omega meson mass in medium can also be obtained by considering the nucleon-antinucleon polarization in medium [125, 156, 157] as in the case of the rho meson. On the other hand, using effective chiral Lagrangians, one finds at finite temperature an even smaller change in the mass of the omega as compared to that of the rho [124]. Because of its small decay width, which is about an order of magnitude smaller than that of rho meson, most dileptons from omega decay in heavy ion collisions are emitted after freeze out, where the medium effects are negligible. However, with appropriate kinematics an omega can be produced at rest in nuclei in reactions induced by the photon, proton and pion. By measuring the decay of the omega into dileptons allows one to determine its properties in the nuclear matter. Such experiments will be carried out at CEBAF [154] and GSI [155].

### 3.3 *The phi meson*

For the phi meson, the situation is less ambiguous than the rho meson as both effective chiral Lagrangian and QCD sum-rule studies predict that its mass decreases in medium. Based on the hidden gauge theory, Song [110] finds that at a temperature of  $T = 200$  MeV the phi meson mass is reduced by about 20 MeV. The main contribution is from the thermal

kaon loop. In QCD sum rules, the reduction is much more appreciable, i.e., about 200 MeV at the same temperature [109]. The latter is due to the significant decrease of the strange quark condensate at finite temperature as a result of the abundant strange hadrons in hot matter. Because of the relative large strange quark mass compared to the up and down quark masses, the phi meson mass in QCD sum rules is mainly determined by the strange quark condensate instead of the four quark condensate as in the case of rho meson mass. However, the temperature dependence of the strange quark condensate in [109] is determined from a non-interacting gas model, so effects due to interactions, which may be important at high temperatures, are not included. In [110], only the lowest kaon loop has been included, so the change of kaon properties in medium is neglected. As shown in [158], this would reduce the phi meson mass if the kaon mass becomes small in medium. Also, vacuum effects in medium due to lambda-antilambda polarization has also been shown to reduce the phi meson in-medium mass [159].

QCD sum rules have also been used in studying phi meson mass at finite density [133], and it is found to decrease by about 25 MeV at normal nuclear matter density.

Current experimental data on phi meson mass from measuring the  $K\bar{K}$  invariance mass spectra in heavy ion collisions at AGS energies do not show a change of the phi meson mass [160]. This is not surprising as these kaon-antikaon pairs are from the decay of phi mesons at freeze out when their properties are the same as in free space. If a phi meson decays in medium, the resulting kaon and antikaon would interact with nucleons, so their invariant mass is modified and can no longer be used to reconstruct the phi meson. However, future experiments on measuring dileptons from photon-nucleus reactions at CEBAF [154] and heavy ion collisions at GSI [155] will provide useful information on the phi meson properties in dense nuclear matter. For heavy ion collisions at RHIC energies, matter with low-baryon chemical potential is expected to be formed in central rapidities. Based on the QCD-sum rule prediction for the mass shift of the phi meson it has been suggested that a low mass phi peak at  $\sim 880$  MeV besides the normal one at 1.02 GeV appears in the dilepton spectrum if a first-order phase transition or a slow cross-over between the quark-gluon plasma and the hadronic matter occurs in the collisions [161, 162]. The low-mass phi peak is due to the nonnegligible duration time for the system to stay near the transition temperature compared with the lifetime of a phi meson in vacuum, so the contribution to dileptons from phi meson decays in the mixed phase becomes comparable to that from their decays at freeze

out. Without the formation of the quark-gluon plasma, the low-mass phi peak is reduced to a shoulder in the dilepton spectrum. Thus, one can use this double phi peaks in the dilepton spectrum as a signature for identifying the quark-gluon plasma to hadronic matter phase transition in ultrarelativistic heavy ion collisions.

## 4 BARYONS

As far as the in medium properties of the baryons are concerned most is known about the nucleon. But also the properties of the hyperons such as the  $\Lambda$  can be determined in the medium by studying hypernuclei. Also the  $\Lambda(1405)$ , which we have discussed in connection with the  $K^-$  optical potential, is another example for in-medium effects of baryons. In the following we will limit ourself on a brief discussion of how the nucleon properties can be viewed in the context of chiral symmetry.

As mentioned briefly in the introduction, medium effects on a nucleon due to many-body interactions have long been studied, leading to an effective mass, which is generally reduced as a result of the finite range of the nucleon interaction and higher order effects. Its relation to chiral symmetry restoration is best seen through the QCD sum rules [2, 3, 163, 164, 165, 166, 167]. Using in-medium condensates, it has been shown that the change of scalar quark condensate in medium leads to an attractive scalar potential which reduces the nucleon mass, while the change of vector quark condensate leads to a repulsive vector potential which shifts its energy. The nucleon scalar and vector self-energies in this study are given, respectively, by

$$\begin{aligned}\Sigma_S &\approx -\frac{8\pi^2}{M_B^2}(\langle\bar{q}q\rangle_\rho - \langle\bar{q}q\rangle_0) \approx -\frac{8\pi^2}{M_B^2}\frac{\Sigma_{\pi N}}{m_u + m_d}\rho_N, \\ \Sigma_V &\approx \frac{64\pi^2}{3M_B^2}\langle q^+q\rangle_\rho = \frac{32\pi^2}{M_B^2}\rho_N,\end{aligned}\tag{19}$$

where the Borel mass  $M_B$  is an arbitrary parameter. With  $M_B \approx m_N$ , and  $m_u + m_d \approx 11$  MeV, these self-energies have magnitude of a few hundred MeV at normal nuclear density. These values are similar to those determined from both the Walecka model [168] and the Dirac-Brueckner-Hartree-Fock (DBHF) approach [169] based on the meson-exchange nucleon-nucleon interaction. Experimental evidences for these strong scalar and vector potentials have been inferred from the proton-nucleus scattering at intermediate energies [170] via the Dirac phe-

nomenology [171, 172, 173, 174] in which the Dirac equation with scalar and vector potentials is solved.

The importance of medium effects on the nucleon mass can be seen from the significant decrease of the  $Q$ -value for the reaction  $NN \rightarrow NNN\bar{N}$  in nuclear medium as a result of the attractive scalar potential. This effect has been included in a number of studies based on transport models. In Fig. 8, theoretical results from these calculations for the antiproton differential cross section in Ni+Ni collisions at 1.85 GeV/nucleon [175, 176, 177] and C+Cu collisions at 3.65 GeV/nucleon [178] are compared with the experimental data from GSI [179] and Dubna [180]. In [175, 176, 177], the relativistic transport model has been used with the antiproton mean-field potential obtained from the Walecka model. The latter has a value in the range of -150 to -250 MeV at normal nuclear matter density. The results of ref. [178] are based on the nonrelativistic Quantum Molecular Dynamics with the produced nucleon and antinucleon masses taken from the Nambu–Jona-Lasinio model. Within this framework, these studies thus show that in order to describe the antiproton data from heavy-ion collisions at subthreshold energies it is necessary to include the reduction of both nucleon and antinucleon masses in nuclear medium. Even at AGS energies, which are above the antiproton production threshold in NN interaction, the medium effects on antiproton may still be important [181]. Indeed, a recent study using the Relativistic Quantum Molecular Dynamics shows that medium modifications of the antiproton properties are important for a quantitative description of the experimental data [182]. We note, however, a better understanding of antiproton annihilation is needed, as only about 5-10% of produced antiprotons can survive, in order to determine more precisely the antiproton in-medium mass from heavy ion collisions.

## 5 RESULTS FROM LATTICE QCD CALCULATIONS

Another source of information about possible in medium changes are lattice QCD-calculations. For a review see e.g. [1]. Presently lattice calculations can only explore systems at finite temperature but vanishing baryon chemical potential. Therefore, their results are not directly applicable to present heavy ion experiments, where system at finite baryon chemical potential are created. But future experiments at RHIC and LHC might succeed in generating a baryon free region. But



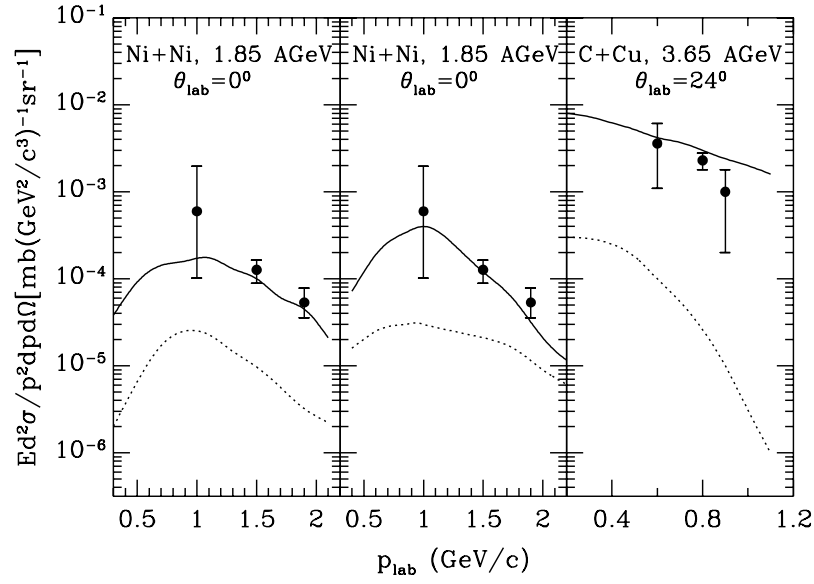


Figure 8: Antiproton momentum spectra from Ni+Ni collisions at 1.85 GeV/nucleon, and C+Cu collisions at 3.85 GeV/nucleon. The left, middle, and right panels are from Refs. [175], [177], and [178], respectively. The solid and dotted curves are from transport model calculations with free and dropping antiproton mass, respectively. The experimental data from ref. [179] for Ni+Ni collisions and from ref. [180] for C+Cu collisions are shown by solid circles.

aside from the experimental aspects, lattice results provide an important additional source where our model understanding can be tested. Lattice calculation are usually carried out in Euclidean space, i.e. in a space with imaginary time. As a consequence plane waves in Minkowski-space translate into decaying exponentials in Euclidean space. At zero temperature, masses of the low lying hadrons are extracted from two point functions, which carry the quantum numbers of the hadron under consideration.

$$C(\tau) = \int d^3x \langle J(x, \tau) J(0) \rangle \sim \sum_i \alpha_i^2 \cdot \exp(-m_i \tau) \quad (20)$$

Here, the sum goes over all hadronic states which carry the quantum numbers of the operator  $J$ . At large imaginary times  $\tau$  only the state with the lowest mass survives, and its mass can be determined from the exponential slope. Zero temperature lattice calculation by now reproduce the hadronic spectrum to a remarkable accuracy [183]. If one wants to extract masses at finite temperature, however, things become more complicated. The reason is that finite temperature on the lattice is equivalent to requiring periodic or anti-periodic boundary conditions in the imaginary time direction for bosons or fermions, respectively. Therefore, one cannot study the correlation functions at arbitrary large  $\tau$  but is restricted to  $\tau_{max} = 1/T$ , where  $T$  is the temperature under consideration. As a result, the lowest lying states cannot be projected out so easily. This has led people to study the correlation functions as a function of the spatial distance, where no restriction to the spatial extent exist, to extract so-called screening masses. At zero temperature the Euclidean time and space direction are equivalent and the screening masses are identical to the actual hadron masses. By analyzing the correlation function in the spatial direction one essentially measures the range of a virtual particle emission. As already demonstrated by Yukawa, who deduced the mass of the pion from the range of the nuclear interaction, this can be used to extract particle masses. Screening masses, however, are not very useful if one wants to study hadronic masses at high temperature, close to  $T_c$ . At these temperatures, the spatial correlators are dominated by the trivial contribution from free independent quarks, and the resulting screening masses simply turn out to be [184]

$$M_{\text{screen}} = n\pi T, \quad (21)$$

where  $n$  is the number of quarks of a given hadron, i.e.  $n = 2$  for mesons and  $n = 3$  for baryons. Only for the pion and sigma meson, lat-

tice calculation found a significant deviation from this simple behavior, which is due to strong residual interactions in these channel. All other hadrons considered so far, such as the  $\rho$  and  $a_1$  as well as the nucleon, exhibit the above screening mass [185, 186]. These trivial contributions from independent quarks are absent, however, in the time-like correlator. There has been one attempt to extract meson masses from the time-like correlators in four flavor QCD [187]. In this case it has been found that the mass of the  $\rho$ -meson does not change significantly below  $T_c$ . Above the critical temperature, on the other hand, the correlation function appeared to be consistent with one of noninteracting quarks, indicating the onset of deconfinement.

There have been attempts to extract the mass of the scalar  $\sigma$  meson as well as that of the pion from so called susceptibilities [188]. These are nothing else than the total four-volume integral over the two-point functions. If the two-point function is dominated by one hadronic pole, then this integral should be inversely proportional to the square of the mass of the lightest hadron under consideration. Using this method, the degeneracy of the pion and sigma mass close to  $T_c$  has been demonstrated.

Finally, let us note that a scenario where all hadron masses scale linearly with the chiral condensate is consistent with strong rise in the energy and entropy density around the phase transition as observed in Lattice calculations [116].

## 6 SUMMARY

In this review, hadron properties, particularly their masses, in the nuclear medium have been discussed. Both conventional many-body interactions and genuine vacuum effects due to chiral symmetry restoration have effects on the hadron in-medium properties.

For Goldstone bosons, which are directly linked to the spontaneously broken chiral symmetry, we have discussed the pion, kaons, and etas. For the pion, its mass is only slightly shifted in the medium due to the small s-wave interactions as a result of chiral symmetry. On the other hand, the strong attractive p-wave interactions due to the  $\Delta(1230)$  and the  $\rho(770)$  lead to a softening of the pion dispersion relation in medium. These effects result from many-body interactions rather from chiral symmetry. Phenomenologically, the small change in the pion in-medium mass does not seem to have any observable effects. However, the effects due to the softened pion in-medium dispersion relation may

be detectable through the enhanced low transverse momentum pions in heavy ion collisions.

For kaons, medium effects due to the Weinberg-Tomozawa vector type interaction, which are negligible for pions in asymmetric nuclear matter, are important as they have only one light quark. However, because of the large explicit symmetry breaking due to the finite strange quark mass higher order corrections in the chiral expansion are non-negligible, leading to an appreciable attractive scalar interaction for both the kaon and the antikaon in the medium. Available experimental data on both subthreshold kaon production and kaon flow are consistent with the presence of this attractive scalar interaction.

For etas, their properties are more related to the explicit breaking of the  $U_A(1)$  axial symmetry in QCD. If the  $U_A(1)$  symmetry is restored in medium as indicated by the instanton liquid model, then their masses are expected to decrease as well. Heavy ion experiments at the CERN-SPS, however, rule out a vast enhancement of the final-state eta yield.

In the case of vector mesons, we have discussed the rho, omega, and phi. At finite temperature, all model calculations agree with the current algebra result that the mass of the rho does not change to order  $T^2$ . To higher order in the temperature and at finite density QCD sum rules predict a dropping of the masses in the medium. Predictions from chiral models, on the other hand, tend to predict an increase of the rho mass with temperature. However, presently available dilepton data from CERN SPS heavy ion experiments are best described assuming a dropping rho mass in dense matter but the errors in the present experimental data are too large to definitely exclude some more conventional explanations.

For baryons, particularly the nucleon, the change of their properties in the nuclear medium has been well-known in studies based on conventional many-body theory. On the other hand, the nucleon mass is also found to be reduced in the medium as a result of the scalar attractive potential related to the quark condensate. A clear separation of the vacuum effects due to the condensate from those of many-body interactions is a topic of current interest and has not yet been resolved. A dropping nucleon effective mass in medium seems to be required to explain the large enhancement of antiproton production in heavy ion collisions at subthreshold energies.

In principle lattice QCD calculations could help answer quite a few of these questions, although they are restricted to systems at finite temperature and vanishing baryon density. However, with the presently available computing power reliable quantitative predictions about in-

medium properties of hadrons are still not available.

The study of in-medium properties of hadrons is a very active field, both theoretically and experimentally. While many questions are still open and require additional measurements and more careful calculations, it is undisputed that the study of in medium properties of hadrons provides us with an unique opportunity to further our understanding about the long range, nonperturbative aspects of QCD.

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## References

- [1] DeTar CE. 1995. *Quark Gluon Plasma 2* ed. R. Hwa, Singapore: World Scientific
- [2] Drukarev EG, Levin EM. 1990. *Nucl. Phys. A* 511:679-700
- [3] Cohen TD, Furnstahl RJ, Griegel DK. 1992. *Phys. Rev. C* 45:1881-93
- [4] Li GQ, Ko CM. 1994. *Phys. Lett. B* 338:118-22
- [5] Brockmann R, Weise W. 1996. *Phys. Lett. B* 367:40-44.
- [6] Schäfer T, Shuryak EV. 1996. Instantons in QCD. *hep-ph/9610451*, to appear in *Rev. Mod. Phys*
- [7] Rho M. 1994. *Phys. Repts.* 240:1-142
- [8] Brown GE, Rho M. 1991. *Phys. Rev. Lett.* 66:2720-3
- [9] Aichelin J, Ko CM. 1985. *Phys. Rev. Lett.* 55:2661-4
- [10] Batko G, Cassing W, Mosel U, Niita K, Wolf G. 1991. *Phys. Lett. B* 256:331-6
- [11] Jeukenne JP, Lejeune A, and Mahaux C. 1976 *Phys. Rep.* 25:83-174
- [12] Mahaux C, Bortignon PF, Broglia RA, Dasso CH. 1985. *Phys. Rep.* 120:1-274

- [13] Brown GE, Weise W, Baym G, Speth J. 1987. *Comments Nucl. Part. Phys.* 17:39-62
- [14] Gelmini G, Ritzi B. 1995. *Phys. Lett. B* 357:431-4
- [15] Brown GE, Rho M. 1996. *Nucl. Phys. A* 596:503-14
- [16] Furnstahl RJ, Tang HB, Serot BD. 1995. *Phys. Rev. C* 52:1368-79
- [17] Stöcker H, Greiner W. 1986. *Phys. Rep.* 137:277-392
- [18] Bertsch GF, Das Gupta S. 1988. *Phys. Rep.* 160:189-233
- [19] Cassing W, Metag V, Mosel U, Niita K. 1990. *Phys. Rep.* 188:363-449
- [20] Aichelin J. 1991. *Phys. Rep.* 202:235-360
- [21] Blättel B., Koch V., and Mosel U. 1993. *Rep. Prog. Phys.* 56:1-62
- [22] Bonasera A, Gulminelli G, Molitoris J. 1994. *Phys. Rep.* 243:1-124
- [23] Ko CM, Li GQ. 1996. *J. Phys. G* 22:1673-725
- [24] Effenberger M, Hombach A, Teis S, Mosel M. 1996. *nucl-th/9607005*
- [25] Effenberger M, Hombach A, Teis S, Mosel M. 1996. *nucl-th/9610022*
- [26] Li BA, Bauer W, Ko CM. *Phys. Lett. B* 382:337-42
- [27] Matsui T, Satz H. 1986. *Phys. Lett. B* 178:416-422
- [28] Gerschel C, Hüfner J. 1992. *Z. Phys. C* 56:171-4
- [29] Gavin S, Vogt R. 1990. *Nucl. Phys. B* 345:104-24
- [30] Cassing W, Ko CM. 1996. *Phys. Lett. B*. In press.
- [31] 1996. *Proc. Quark Matter 96, Nucl. Phys. A610*, ed. P. Braun-Munzinger *et al.*
- [32] Kharzeev D, Satz H. 1995. In *Quark Gluon Plasma 2*, ed. R. Hwa, Singapore: World Scientific.
- [33] Ericson T, Weise W. 1988. *Pions and Nuclei*. Oxford: Clarendon Press
- [34] Delorme J, Ericson M, Ericson TEO. 1992. *Phys. Lett. B* 291:379-84
- [35] Thorsson V, Wirzba A. 1995. *Nucl. Phys. A* 589:633-48
- [36] Lutz M, Klimt S., Weise W. 1992. *Nucl. Phys. A* 542 521-58
- [37] Weinberg S. 1966. *Phys. Rev. Lett.* 17:616-21
- [38] Bertsch GF, Brown GE, Koch V, Li BA. 1988. *Nucl. Phys. A* 490:745-55
- [39] Xiong L, Ko CM, Koch V. 1993. *Phys. Rev. C* 47:788-94
- [40] Friedman B, Pandharipande VR, Usmani QN. 1981. *Nucl. Phys. A*, 372:483-95
- [41] Ericson M, Ericson TEO. 1966. *Ann. Phys. (N.Y.)* 36:323-362

- [42] Brown GE, Oset E, Vacas MV, Weise W. 1989. *Nucl. Phys. A* 505:823-34
- [43] Ko CM, Xia LH, Siemens PJ. 1989. *Phys. Lett. B* 231:16-20
- [44] Xia LH, Siemens PJ, Soyeur M. 1994. *Nucl. Phys. A* 578:493-510
- [45] Henning PA, Umezawa H. 1994. *Nucl. Phys. A* 571:617-644
- [46] Korpa CL, Malfliet R. 1995. *Phys. Rev. C* 52:2756-61
- [47] Helgeson J, Randrup J. 1995. *Ann. Phys. (NY)* 244:12-66
- [48] Hennio T, et al. 1993. *Phys. Lett. B* 303:236-39
- [49] C. Garde, private communication.
- [50] Gale C, Kapusta J. 1987. *Phys. Rev. C* 35:2107-2116
- [51] Korpa CL, Pratt S. 1990. *Phys. Rev. Lett.* 64:1502-5
- [52] Korpa CL, Xiong L, Ko CM, Siemens PJ. 1990. *Phys. Lett. B* 246:333-6
- [53] Roche G, et al. 1988. *Phys. Rev. Lett.* 61:1069-72
- [54] Naudet C, et al. 1988. *Phys. Rev. Lett.* 62:2652-55
- [55] Ehehalt W, Cassing W, Engel A, Mosel U, Wolf G. 1993. *Phys. Lett. B* 298:31-5
- [56] Zipprich J, Fuchs C, Lehmann E, Sehn L, Huang S W, Faessler A. 1997. *J. Phys. G* 23:L1-L6
- [57] Odyniec G. 1987. *Proc. 8th High Energy Heavy Ion Study*, ed. JW Harris, GJ Wozniak, pp. 215. Berkeley
- [58] Berg FD, et al. 1990. *Z.Phys. A* 340:297-302
- [59] Gerber P, Leutwyler H. 1989. *Nucl. Phys. B* 321:387-429
- [60] Shuryak E. 1990. *Phys. Rev. D* 42:1764-76
- [61] Shuryak E. 1991. *Nucl. Phys. A* 533:761-788
- [62] Koch V, Bertsch GF. 1993. *Nucl. Phys. A* 552:591-604
- [63] Song C. 1994. *Phys. Rev. D* 49:1556-65
- [64] Song C, Koch V, Lee SH, Ko CM. 1996. *Phys. Lett. B* 366:379-84
- [65] Koch V, Song C. 1996. *Phys. Rev. C* 54:1903-13
- [66] Song C, Koch V. 1996. *Phys. Rev. C* 54:3218-31
- [67] Stroebelen H, et al. (NA35 collaboration). 1988. *Z.Phys. C* 38:89-96
- [68] Stroebelen H, et al. (NA35 collaboration). 1990. *Nucl. Phys. A* 525:59c-66c
- [69] Kaplan DB, Nelson AE. 1986. *Phys. Lett. B* 175:57-63
- [70] Nelson AE, Kaplan DB. 1987. *Phys. Lett. B* 192:193-7
- [71] Politzer HD, Wise MB. 1991. *Phys. Lett. B* 273:156-62
- [72] Brown GE, Lee CH, Rho M, Thorsson V. 1994. *Nucl. Phys. A* 567:937-56
- [73] Jenkins E. 1992. *Nucl. Phys. B* 368:190-203
- [74] Martin, A.D. 1981. *Nucl. Phys. B* 179:33-48

- [75] Dalitz RH, Wong TC, Rajasekaran G. 1967. *Phys. Rev* 153:1617-23
- [76] Lee CH, Brown GE, Rho M. 1994. *Phys. Lett. B* 335:266-72
- [77] Kaiser N, Siegel PB, Weise W. 1995. *Nucl. Phys. A* 594:325-45
- [78] Koch V. 1994. *Phys. Lett. B* 337:7-11
- [79] Friedman E, Gal A, Batty CJ. 1993. *Phys. Lett. B* 308:6-10
- [80] Waas T, Kaiser N, Weise W. 1996. *Phys. Lett. B* 365:12-6
- [81] Brown GE, Thorsson V, Kubodera K, Rho M. 1992. *Phys. Lett. B* 291:355-62
- [82] Brown GE, Bethe HA. 1994. *Astrophys. Jour.* 423:659-64
- [83] Fang XS, Ko CM, Li GQ, Zheng YM. 1994. *Phys. Rev. C* 49:R608-11
- [84] Fang XS, Ko CM, Li GQ, Zheng YM. 1994. *Nucl. Phys. A* 575:766-90
- [85] Li GQ, Ko CM. 1995. *Phys. Lett. B* 349:405-10
- [86] Li GQ, Ko CM, Li BA. 1995. *Phys. Rev. Lett.* 74:235-8
- [87] Li GQ, Ko CM. 1995. *Nucl. Phys. A* 594:460-82
- [88] Miskowiec D *et al.*, *Phys. Rev. Lett.* 72:3650-3
- [89] Ritman J, the FOPI collaboration. 1995. *Z.Phys. A* 352:355-7
- [90] David C, Hartnack C, Kerveno M, Le Pallec JC, Aichelin J. 1996. *Nucl. Phys. A*. Submitted
- [91] Maruyama T, Cassing W, Mosel U, Teis S, Weber K. 1994. *Nucl. Phys. A* 573:653-75
- [92] Gibson BF, Hungerford EV. 1995 *Phys. Rep.* 257:349
- [93] Chen CM, Ernst DJ. 1992. *Phys. Rev. C* 45:2019-22
- [94] Brown GE, Dover CB, Siegel PB, Weise W. 1988. *Phys. Rev. Lett.* 60:2723-26
- [95] Li GQ, Ko CM, Fang XS. 1994. *Phys. Lett. B* 329:149-56
- [96] Li GQ, Ko CM. 1996. *Phys. Rev. C* 54:R2159-R2162
- [97] Schröter A, *et al.* 1994. *Z.Phys. A* 350:101-13
- [98] Ritman. 1996. private communications
- [99] Shuryak E. 1994. *Comm. Nucl. Part. Phys.* 21:235-248
- [100] Schäfer T. 1996. *Phys. Lett. B* 389:455-51
- [101] Kapusta J, Kharzeev D, McLerran L. 1996. *Phys. Rev. D* 53:5028-33
- [102] Huang Z, Wang XN. 1996. *Phys. Rev. D* 53:5034-41
- [103] Berg FD, *et al.* 1994. *Phys. Rev. Lett.* 72:977-80
- [104] Drees A. 1996. *Phys. Lett. B* 388:380-3
- [105] Koch V. 1990. *Proc. Pittsburgh Workshop on Soft Lepton Pair and Photon Production*, ed. JA Thompson, p. 251. New York: Nova
- [106] Wolf G, Cassing W, Mosel U. 1993. *Nucl. Phys. A* 552:549-70
- [107] Li GQ, Ko CM. 1995. *Nucl. Phys. A* 583:731-48



- [108] Weinberg S. 1967. *Phys. Rev. Lett.* 18:507-9
- [109] Asakawa M, Ko CM. 1994. *Nucl. Phys. A* 572:732-48
- [110] Song C. 1996. *Phys. Lett. B* 388:141-6
- [111] Adami C, Brown GE. 1993. *Phys. Rep.* 234:1-71
- [112] Friman B, Rho M. 1996. *Nucl. Phys. A* 606:303-19
- [113] Brown GE, Buballa M, Rho M. 1996. *Nucl. Phys. A* 609:519-36
- [114] Donoghue J, Golowich E, Holstein B. 1992. *Dynamics of the Standard Model*. Cambridge: Cambridge University Press, Cambridge, U.K.
- [115] Pisarski R. 1982. *Phys. Lett. B* 110:155-59
- [116] Koch V, Brown GE. 1993. *Nucl. Phys. A* 560:345-64
- [117] Sakurai JJ. 1969. *Currents and mesons*. Chicago: University of Chicago Press, Chicago
- [118] Gale C, Kapusta JJ. 1991. *Nucl. Phys. B* 357:65-89
- [119] Reinders LJ, Rubinstein H, Yazaki S. 1985. *Phys. Rep.* 127:1-97
- [120] Kapusta JJ, Shuryak EV. 1994. *Phys. Rev. D* 49:4694-704
- [121] Dey M, Eletsky VL, Ioffe BL. 1990. *Phys. Lett. B* 252:620-4
- [122] Eletsky VL, Ioffe BL. 1995. *Phys. Rev. D* 51:2371-6
- [123] Pisarski RD. 1995. *Phys. Rev. D* 52:R3773-R3776
- [124] Song C. 1993. *Phys. Rev. D* 48:1375-89
- [125] Song CS, Xia PW, Ko CM. 1995 *Phys. Rev. C* 52:408-11
- [126] Rapp R, Wambach J. 1993. *Phys. Lett. B* 315:220-5
- [127] Chanfray G, Schuck P. 1993. *Nucl. Phys. A* 555:329-54
- [128] Herrmann M, Friman B, Nörenberg W. 1993. *Nucl. Phys. A* 560:411-36
- [129] Asakawa M, Ko CM, Levai P, Qiu XJ. *Phys. Rev. C* 46:R1159-R1162
- [130] Klingl F, Weise W. 1996. *Nucl. Phys. A* 606:329-38
- [131] Klingl F, Weise W. 1997. Proceedings of the Intl. Workshop XXV on Gross Properties of Nuclei and Nuclear Excitations ed. H. Feldmeier and W. Nörenberg, Hirschegg, Austria, 1997.
- [132] Klingl F, Weise W. in preparation
- [133] Hatsuda T, Lee SH. 1992. *Phys. Rev. C* 46:R34-R38
- [134] Asakawa M, Ko CM. 1993. *Nucl. Phys. A* 560:399-410
- [135] Friman B, Soyeur M. 1996. *Nucl. Phys. A* 600:477-90
- [136] Eletsky VL, Ioffe BL. 1996. *hep-ph/9609229*
- [137] Friman B, Pirner H. 1997 *nucl-th/9701016*
- [138] Xiong L, Wu ZG, Ko CM, Wu JQ. 1990. *Nucl. Phys. A* 512:772-86
- [139] Wolf G, Batko G, Cassing W, Mosel U, Niita K, Schäfer M. 1990. *Nucl. Phys. A* 517:615-38

- [140] Porter R.J. et al. 1997, LBNL-Report LBNL-39957, submitted to Phys. Rev. Lett.
- [141] Drees A. 1997 Proc. Quark Matter 96, Nucl. Phys. A610, ed. P. Braun-Munzinger
- [142] Agakichiev G, *et al.* 1995. *Phys. Rev. Lett.* 75:1272-5
- [143] Masera M for the HELIOS Collaboration. 1995. *Nucl. Phys. A* 590:93c-102c
- [144] Cassing W, Eehalt W, Ko CM. 1995. *Phys. Lett. B* 363:35-40
- [145] Srivastava DK, Sinha B, Gale C. *Phys. Rev. C* 53:R567-R571
- [146] Haglin K. 1996. *Phys. Rev. C* 53:R2606-R2609
- [147] Baier R, Dirks M, Redlich R. 1996. *help-ph/9610210*
- [148] Chanfray G, Rapp R, Wambach J. 1996. *Phys. Rev. Lett.* 76:368-71
- [149] LI GQ, Ko CM, Brown GE. 1995. *Phys. Rev. Lett.* 75:4007-10
- [150] LI GQ, Ko CM, Brown GE. 1996. *Nucl. Phys. A* 606:568-606
- [151] LI GQ, Ko CM, Brown GE, Sorge H. 1996. *Nucl. Phys. A* 611:539-67
- [152] Song CS, Lee SH, Ko CM. 1995. *Phys. Rev. C* 52:R476-9
- [153] Kharzeev, Satz H. 1994. *Phys. Lett. B* 340:167-70
- [154] Freedom BM. 1995. *Proc. Internal. Workshop XXXIII on Gross Properties of Nuclei and Nuclear Excitations*, pp. 273-82. GSI
- [155] Stroth J. 1995. *Proc. Internal. Workshop XXXIII on Gross Properties of Nuclei and Nuclear Excitations*, pp. 202-17. GSI
- [156] Jean HC, Piekarewicz J, Williams AG. 1994. *Phys. Rev. C* 49:1981-8
- [157] Shiomo H, Hatsuda T. 1994. *Phys. Lett. B* 334:281-6
- [158] Ko CM, Lévai P, Qiu XJ. 1992. *Phys. Rev. C* 45:1400-2
- [159] Kuwabara H, Hatsuda T. 1995. *Prog. Theo. Phys.* 94:1163-7
- [160] Akiba Y, *et al.* 1996. *Phys. Rev. Lett.* 76:2021-24
- [161] Asakawa M, Ko CM. 1994. *Phys. Lett. B* 322:33-7
- [162] Asakawa M, Ko CM. 1994. *Phys. Rev. C* 50:3064-8
- [163] Cohen TD, Furnstahl RJ, Griegel DK. 1991. *Phys. Rev. Lett.* 67:961-4
- [164] Jin XM, Cohen TD, Furnstahl RJ, Griegel DK. 1993. *Phys. Rev. C* 47:2882-900
- [165] Jin XM, Nielsen M, Cohen TD, Furnstahl RJ, Griegel DK. 1994. *Phys. Rev. C* 49:464-77
- [166] Hatsuda T, Higaasen T, Prakash M. 1990. *Phys. Rev. C* 42:2212-21
- [167] Adami C, Brown GE. 1990 *Z. Phys. A* 340:93-100
- [168] Serot BD, Walecka JD. 1986. *Adv. Nucl. Phys.* 16:1
- [169] Machleidt R. 1989. *Adv. Nucl. Phys.* 19:189-376

- [170] Ray L, Hoffmann GW, Coker WR. 1992. *Phys. Rep.* 212:223-328
- [171] Arnold LG, Clark BC, Mercer RL. 1979. *Phys. Rev. C* 19:917-22
- [172] Kobbs AM, Cooper ED, Johansson JI, Sherif HS. 1985. *Nucl. Phys. A* 445:605-24
- [173] Hama S, Clark BC, Cooper ED, Sherif HS, Mercer RL. 1990. *Phys. Rev. C* 41:2737-55
- [174] Cooper ED, Hama S, Clark BC, Mercer RL. 1993. *Phys. Rev. C* 47:297-311
- [175] Li GQ, Ko CM. 1994. *Phys. Rev. C* 50:1725-8
- [176] Li GQ, Ko CM, Fang XS, Zheng YM. 1994. *Phys. Rev. C* 49:1139-48
- [177] Teis S, Cassing W, Maruyama T, Mosel U. 1994. *Phys. Rev. C* 50:388-405
- [178] Batko G, Faessler A, Huang SW, Lehmann E, Puri RK. 1994. *J. Phys. G* 20:461-5
- [179] Schröter A, *et al.* 1994. *Nucl. Phys. A* 553:775c-778c
- [180] Baldin AA, *et al.* 1992. *Nucl. Phys. A* 519:407c-411c
- [181] Koch V, Ko CM, Brown GE. 1991. *Phys. Lett. B* 265:29-34
- [182] Spieles C, Bleicher M, Jahns A, Mattiello R, Sorge H, Stöcker H, Greiner W. 1996. *Phys. Rev. C* 53:2011-3
- [183] Sexton J, Vaccarino A, Weingarten D, Butler F, Chen H. 1994. *Nucl. Phys. B* 430:179-228
- [184] Eletskii VL, Ioffe BL. 1988. *Sov. J. Nucl. Phys* 48:602-
- [185] DeTar CE, Kogut J. 1987. *Phys. Rev. D* 36:2828-39
- [186] Born KD, Gupta S, Irbäck A, Karsch F, Laermann E, Petersson B, Satz H, (MT<sub>C</sub> Collaboration). 1991. *Phys. Rev. Lett.* 67:302-5
- [187] Boyd G, Gupta S, Karsch F, Laermann E, Peterson B, Redlich K. 1995. *Phys. Lett. B* 349:170-6
- [188] Learmann E, 1996 *Proc. Quark Matter 96, Heidelberg, to appear in Nucl. Phys. A*, ed. J Specht. North Holland

# PROPERTIES OF HADRONS IN THE NUCLEAR MEDIUM

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## Abstract

This review is devoted to the discussion of hadron properties in the nuclear medium and its relation to the partial restoration of chiral symmetry. Special attention is given to disentangle in-medium effects due to conventional many-body interactions from those due to the change of the chiral condensate. In particular, we shall discuss medium effects on the Goldstone bosons (pion, kaon and eta), the vector mesons (rho, omega, phi), and the nucleon. Also, for each proposed in-medium effect the experimental consequence and results will be reviewed.

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## 1 INTRODUCTION

The atomic nucleus provides a unique laboratory to study the long range and bulk properties of QCD. Whereas QCD is well tested in the perturbative regime rather little is known about its properties in the long range, nonperturbative region. One of the central nonperturbative properties of QCD is the spontaneous breaking chiral symmetry in the ground state resulting in a nonvanishing scalar quark condensate,  $\langle \bar{q}q \rangle \neq 0$ . It is believed and supported by lattice QCD calculations [1] that at temperatures around 150 MeV, QCD undergoes a phase transition to a chirally restored phase, characterized by the vanishing of the order parameter, the chiral condensate  $\langle \bar{q}q \rangle$ . This is supported by results obtained within chiral perturbation theory [2, 3]. Effective chiral models predict that a similar transition also takes place at finite nuclear density.

The only way to create macroscopic, strongly interacting systems at finite temperature and/or density in the laboratory is by colliding heavy nuclei at high energies. Experiments carried out at various bombarding energies, ranging from 1 AGeV (BEVALAC, SIS) to 200 AGeV (SPS), have established that one can generate systems of large density but moderate temperatures (SIS, BEVALAC), systems of both large density and temperature (AGS) as well as systems of low density and high temperatures (SPS). Therefore, a large region of the QCD phase diagram can be investigated through the variation of the bombarding energy. But in addition the atomic nucleus itself represents a system at zero temperature and finite density. At nuclear density the quark condensate is estimated to be reduced by about 30% [2, 3, 4, 5] so that effects due to the change of the chiral order parameter may be measurable in reactions induced by a pion, proton or photon on the nucleus.

Calculations within the instanton liquid model [6] as well as results from phenomenological models for hadrons [7] suggest that the properties of the light hadrons, such as masses and couplings, are controlled by chiral symmetry and its spontaneous as well as explicit breaking. Confinement seems to play a lesser role. If this is the correct picture of the low energy excitation of QCD, hadronic properties should depend on the value of the chiral condensate  $\langle \bar{q}q \rangle$ . Consequently, we should expect that the properties of hadrons change considerably in the nuclear environment, where the chiral condensate is reduced. Indeed, based on the restoration of scale invariance of QCD, Brown and Rho have argued that masses of nonstrange hadrons would scale with the quark condensate and thus decrease in the nuclear medium [8]. This has since stimulated extensive theoretical and experimental studies on hadron in-medium properties.

By studying medium effects on hadronic properties one can directly test our understanding of those non-perturbative aspects of QCD, which are responsible for the light hadronic states. The best way to investigate the change of hadronic properties in experiment is to study the production of particles, preferably photons and dileptons, as they are not affected by final-state interactions. Furthermore, since vector mesons decay directly into dileptons, a change of their mass can be seen directly in the dilepton invariant mass spectrum. In addition, as we shall discuss, the measurement of subthreshold particle production such as kaons [9] and antiprotons [10] may also reveal some rather interesting in-medium effects.

Of course a nucleus or a hadronic system created in relativistic heavy ion collisions are strongly interacting. Therefore, many-body excitations can carry the same quantum numbers as the hadrons under consideration and thus can mix with the hadronic states. In addition, in the nuclear environment ‘simple’ many-body effects such as the Pauli principle are at work, which, as we shall discuss, lead to considerable modifications of hadronic properties in some cases. How these effects are related to the partial restoration of chiral symmetry is a new and unsolved question in nuclear many-body physics. One example is the effective mass of a nucleon in the medium, which was introduced long time ago [11, 12] to model the momentum dependence of the nuclear force, which is due to its finite range, or to model its energy dependence, which results from higher order (2-particle - 1-hole etc.) corrections to the nucleon self energy. On the other hand, in the relativistic mean-field description one also arrives at a reduced effective mass of the nucleon. According to [13], this is due to so-called virtual pair corrections and may be

related to chiral symmetry restoration as suggested by recent studies [3, 14, 15, 16].

Another environment, which is somewhat ‘cleaner’ from the theorists point of view, is a system at finite temperature and vanishing baryon chemical potential. At low temperatures such a system can be systematically explored within the framework of chiral perturbation theory, and essentially model-independent statements about the effects of chiral restoration on hadronic masses and couplings may thus be given. In the high temperature regime eventually lattice QCD calculations should be able to tell us about the properties of hadrons close to the chiral transition temperature. Unfortunately, such a system is very difficult to create in the laboratory.

Since relativistic heavy ion collisions are very complex processes, one has to resort to careful modeling in order to extract the desired in-medium correction from the available experimental data. Computationally the description of these collisions is best carried out in the framework of transport theory [17, 18, 19, 20, 21, 22, 23] since nonequilibrium effects have been found to be important at least at energies up to 2 AGeV. At higher energies there seems to be a chance that one can separate the reaction in an initial hard scattering phase and a final equilibrated phase. It appears that most observables are dominated by the second stage of the reaction so that one can work with the assumption of local thermal equilibrium. However, as compared to hydrodynamical calculations, also here the transport approach has some advantages since the freeze out conditions are determined from the calculation and are not needed as input parameter. As far as a photon or a proton induced reaction on nuclei is concerned, one may resort to standard reaction theory. But recent calculations seem to indicate that these reaction can also be successfully treated in the transport approach [24, 25, 26]. Thus it appears possible that one can explore the entire range of experiments within one and the same theoretical framework which of course has the advantage that one reduces the ambiguities of the model to a large extent by exploring different observables.

This review is devoted to the discussion of in-medium effects in the hadronic phase and its relation to the partial restoration of chiral symmetry. We, therefore, will not discuss another possible in-medium effect which is related to the deconfinement in the Quark Gluon Plasma, namely the suppression of the  $J/\Psi$ . This idea, which has been first proposed by Matsui and Satz [27] is based on the observation that due to the screening of the color interaction in the Quark Gluon Plasma, the  $J/\Psi$  is not bound anymore. As a result, if such a Quark Gluon Plasma

is formed in a relativistic heavy ion collision, the abundance of  $J/\Psi$  should be considerably reduced. Of course also this signal suffers from more conventional backgrounds, namely the dissociation of the  $J/\Psi$  due to hadronic collisions [28, 29, 30]. To which extent present data can be understood in a purely hadronic scenario is extensively debated at the moment [31]. We refer the reader to the literature for further details [32].

In this review we will concentrate on proposed in-medium effects on hadrons. Whenever possible, we will try to disentangle conventional in-medium effects from those we believe are due to new physics, namely the change of the chiral condensate. We first will discuss medium effects on the Goldstone bosons, such as the pion, kaon and eta. Then we will concentrate on the vector mesons, which have the advantage that possible changes in their mass can be directly observed in the dilepton spectrum. We further discuss the effective mass of a nucleon in the nuclear medium. Finally, we will close by summarizing the current status of Lattice QCD calculations concerning hadronic properties. As will become clear from our discussion, progress in this field requires the input from experiment. Many question cannot be settled from theoretical consideration alone. Therefore, we will always try to emphasize the observational aspects for each proposed in-medium correction.

## 2 GOLDSTONE BOSONS

This first part is devoted to in-medium effects of Goldstone bosons, specifically the pion, kaon and eta. The fact that these particles are Goldstone bosons means that their properties are directly linked to the spontaneous breakdown of chiral symmetry. Therefore, rather reliable predictions about their properties can be made using chiral symmetry arguments and techniques, such as chiral perturbation theory. Contrary to naive expectations, the properties of Goldstone bosons are rather robust with respect to changes of the chiral condensate. On second thought, however, this is not so surprising, because as long as chiral symmetry is spontaneously broken there will be Goldstone bosons. The actual nonvanishing value of their mass is due to the explicit symmetry breaking, which – at least in case of the pion – are small. Changes in their mass, therefore, are associated with the sub-leading explicit symmetry breaking terms and are thus small. We should stress, however, that in case of the kaon the symmetry breaking terms are considerably larger, leading to sizeable corrections to their mass at finite density as



we shall discuss below.

## 2.1 *The pion*

Of all hadrons, the pion is probably the one where in-medium corrections are best known and understood. From the theoretical point of view, pion properties can be well determined because the pion is such a ‘good’ Goldstone boson. The explicit symmetry breaking terms in this case are small, as they are associated with the current masses of the light quarks. Therefore, the properties of the pion at finite density and temperature can be calculated using for instance chiral perturbation theory. But more importantly, there exists a considerable amount of experimental data ranging from pionic atoms to pion-nucleus experiments to charge-exchange reactions. These data all address the pion properties at finite density which we will discuss now.

### FINITE DENSITY

The mass of a pion in symmetric nuclear matter is directly related to the real part of the isoscalar-s-wave pion optical potential. This has been measured to a very high precision in pionic atoms [33], and one finds to first order in the nuclear density  $\rho$

$$\Delta m_\pi^2 = -4\pi(b_0)_{\text{eff}}\rho, \quad (b_0)_{\text{eff}} \simeq -0.024m_\pi^{-1}. \quad (1)$$

At nuclear matter density the shift of the pion mass is

$$\frac{\Delta m_\pi}{m_\pi} \sim 8\%, \quad (2)$$

which is small as expected from low energy theorems based on chiral symmetry. Theoretically, the value of  $(b_0)_{\text{eff}}$  is well understood as the combined contribution from a single scattering (impulse approximation) involving the small s-wave isoscalar  $\pi N$  scattering length  $b_0 = -0.010(3)m_\pi^{-1}$  and the contribution from a (density dependent) correlation or re-scattering term, which is dominated by the comparatively large isovector-s-wave scattering length  $b_1 = -0.091(2)m_\pi^{-1}$ , i.e.,

$$(b_0)_{\text{eff}} = b_0 - [1 + \frac{m_\pi}{M_N}(b_0^2 + 2b_1^2)] < \frac{1}{r} >, \quad (3)$$

where the correlation length  $< 1/r > \simeq 3p_f/(2\pi)$  is essentially due to the Pauli exclusion principle. The weak and slightly repulsive s-wave

pion potential has also been found in calculations based on chiral perturbation theory [34, 35] as well as the phenomenological Nambu - Jona-Lasinio model [36]. This is actually not too surprising because both amplitudes,  $b_0$  and  $b_1$ , are controlled by chiral symmetry and its explicit breaking [33, 37].

Pions at finite momenta, on the other hand, interact very strongly with nuclear matter through the p-wave interaction, which is dominated by the  $P_{33}$  delta-resonance. This leads to a strong mixing of the pion with nuclear excitations such as a particle-hole and a delta-hole excitation. As a result the pion-like excitation spectrum develops several branches, which to leading order correspond to the pion, particle-hole, and delta-hole excitations. Because of the attractive p-wave interaction the dispersion relation of the pionic branch becomes considerably softer than that of a free pion. A simple model, which has been used in practical calculations [38, 39] is the so-called delta-hole model [33, 40] which concentrates on the stronger pion-nucleon-delta interaction ignoring the nucleon-hole excitations. In this model, the pion dispersion relation in nuclear medium can be written as

$$\omega(\mathbf{k}, \rho) = m_\pi^2 + \mathbf{k}^2 + \Pi(\omega, \mathbf{k}), \quad (4)$$

where the pion self-energy is given by

$$\Pi(\omega, \mathbf{k}) = \frac{\mathbf{k}^2 \chi(\omega, \mathbf{k})}{1 - g' \chi(\omega, \mathbf{k})}, \quad (5)$$

with  $g' \approx 0.6$  the Migdal parameter which accounts for short-range correlations, also known as the Ericson-Ericson - Lorentz-Lorenz effect [41]. The pion susceptibility  $\chi$  is given by

$$\chi(\omega, \mathbf{k}) \approx \frac{8}{9} \left( \frac{f_{\pi N \Delta}}{m_\pi} \right)^2 \frac{\omega_R}{\omega^2 - \omega_R^2} \exp(-2\mathbf{k}^2/b^2) \rho, \quad (6)$$

where  $f_{\pi N \Delta} \approx 2$  is the pion-nucleon-delta coupling constant,  $b \approx 7m_\pi$  is the range of the form factor, and  $\omega_R \approx \frac{\mathbf{k}^2}{2m_\Delta} + m_\Delta - m_N$ .

The pion dispersion relation obtained in this model is shown in Fig. 1. The pion branch in the lower part of the figure is seen to become softened, while the delta-hole branch in the upper part of the figure is stiffened. Naturally this model oversimplifies things and once couplings to nucleon-hole excitations and corrections to the width of the delta are consistently taken into account, the strength of the delta-hole branch is significantly reduced [42, 43, 44, 45, 46, 47]. However, up to momenta of about  $2m_\pi$  the pionic branch remains pretty narrow and considerably

softer than that of a free pion. The dispersion relation of the pion can be measured directly using  $(^3\text{He}, t)$  charge exchange reactions [48]. In these experiments the production of coherent pion have been inferred from angular correlations between the outgoing pion and the transferred momentum. The corresponding energy transfer is smaller than that of free pions indicating the attraction in the pion branch discussed above. More detailed measurements of this type are currently being analyzed [49].

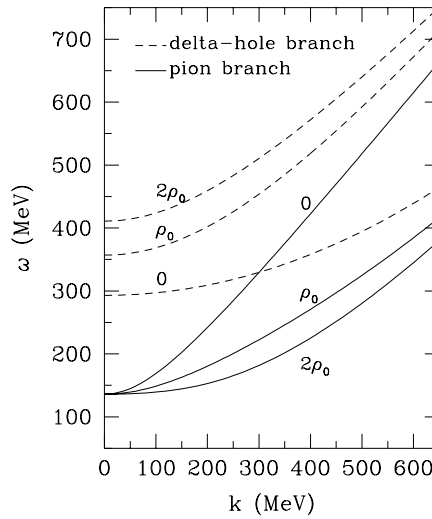


Figure 1: Pion dispersion relation in the nuclear medium. The normal nuclear matter density is denoted by  $\rho_0$ .

In the context of relativistic heavy ion collisions interest in the in-medium modified pion dispersion relation got sparked by the work of Gale and Kapusta [50]. Since the density of states is proportional to the inverse of the group velocity of these collective pion modes, which can actually vanish depending on the strength of the interaction (see Fig. 1), they argued that a softened dispersion relation would lead to a strong enhancement of the dilepton yield. This happens at invariant mass close to twice the pion mass because the number of pions with low energy will be considerably increased. However, as pointed out by Korpa and Pratt [51] and further explored in more detail in [52], the

initially proposed strong enhancement is reduced considerably once corrections due to gauge invariance are properly taken into account. Other places, where the effect of the pion dispersion relation is expected to play a role, are the inelastic nucleon-nucleon cross section [38] and the shape of the pion spectrum [39, 55, 56]. In the latter case one finds [39] that as a result of the attractive interaction the yield of pions at low  $p_t$  is enhanced by about a factor of 2. This low  $p_t$  enhancement is indeed seen in experimental data for  $\pi^-$  production from the BEVALAC [57] and in more recent data from the TAPS collaboration [58], which measures neutral pions. Certainly, in order to be conclusive, more refined calculations are needed, and they are presently being carried out [47].

#### FINITE TEMPERATURE

The pionic properties at finite temperature instead of finite density are qualitatively very similar, although the pion now interacts mostly with other pions instead of nucleons. Again, due to chiral symmetry, the s-wave interaction among pions is small and slightly repulsive, leading to a small mass shift of the pion [59].

And analogous to the case of nuclear matter, there is a strong attractive p-wave interaction, which is now dominated by the  $\rho$ -resonance. This again results in a softened dispersion relation which, however, is not as dramatic as that obtained at finite density [60, 61, 62, 63]. Also the phenomenological consequences are similar. The modified dispersion relation results in an enhancement of dileptons from the pion annihilation by a factor of about two at invariant masses around 300 – 400 MeV [64, 65, 66]. Unfortunately, in the same mass range other channels dominate the dilepton spectrum (see discussion in section 3.1) so that this enhancement cannot be easily observed in experiment [65]. The modified dispersion relation has also been invoked in order to explain the enhancement of low transverse momentum pions observed at CERN-SPS heavy ion collisions [61, 67, 68]. However, at these high energies the expansion velocity of the system, which is mostly made out of pions, is too fast for the attractive pion interaction to affect the pion spectrum at low transverse momenta [62]. This is different at the lower energies around 1 GeV. There, the dominant part of the system and the source for the pion potential are the nucleons, which move considerably slower than pions. Therefore, pions have a chance to leave the potential well before it has disappeared as a result of the expansion.

To summarize this section on pions, the effects for the pions at finite density as well as finite temperature are dominated by p-wave reso-

nances, the  $\Delta(1230)$  and the  $\rho(770)$ , respectively. Both are not related directly to chiral symmetry and its restoration but rather to what we call here many-body effects, which of course does not make them less interesting. The s-wave interaction is small because the pion is such a ‘good’ Goldstone Boson; probably too small to have any phenomenological consequence (at least for heavy ion collisions).

## 2.2 Kaons

Contrary to the pion, the kaon is not such a good Goldstone boson. Effects of the explicit chiral symmetry breaking are considerably bigger, as one can see from the mass of the kaon, which is already half of the typical hadronic mass scale of 1 GeV. In addition, since the kaon carries strangeness, its behavior in non-strange, isospin symmetric matter will be different from that of the pion. Rather interesting phenomenological consequences arise from this difference such as a possible condensation of antikaons in neutron star matter [69, 70, 71, 72].

### CHIRAL LAGRANGIAN

This difference can best be exemplified by studying the leading order effective  $SU(3)_L \times SU(3)_R$  Lagrangian obtained in heavy baryon chiral perturbation theory [72, 73]

$$\begin{aligned} \mathcal{L}_0 = & \frac{f^2}{4} \text{Tr } \partial^\mu U \partial_\mu U^\dagger + \frac{f^2}{2} r \text{Tr } M_q (U + U^\dagger - 2) \\ & + \text{Tr } \bar{B} i v_\mu \mathcal{D}^\mu B + 2D \text{Tr } \bar{B} S^\mu \{A_\mu, B\} \\ & + 2F \text{Tr } \bar{B} S^\mu [A_\mu, B]. \end{aligned} \quad (7)$$

In the above formula, we have  $U = \exp(2i\pi/f)$  with  $\pi$  and  $f$  being the pseudoscalar meson octet and their decay constant, respectively;  $B$  is the baryon octet;  $v_\mu$  is the four velocity of the heavy baryon ( $v^2 = 1$ ); and  $S_\mu$  stands for the spin operator  $S_\mu = \frac{1}{4}\gamma_5[v_\nu\gamma_\nu, \gamma_\mu]$  and  $v_\mu S^\mu = 0$ ;  $M_q$  is the quark mass matrix; and  $r$ ,  $D$  as well as  $F$  are empirically determined constants. Furthermore,

$$\mathcal{D}_\mu B = \partial_\mu B + [V_\mu, B] \quad (8)$$

$$V_\mu = \frac{1}{2}(\xi \partial_\mu \xi^\dagger + \xi^\dagger \partial_\mu \xi), \quad A_\mu = \frac{i}{2}(\xi \partial_\mu \xi^\dagger - \xi^\dagger \partial_\mu \xi), \quad (9)$$

with  $\xi^2 = U$ .

To leading order in the chiral counting the explicit symmetry breaking term  $\sim \text{Tr } M_q (U + U^\dagger - 2)$  gives rise to the masses of the Goldstone

bosons. The interesting difference between the behavior of pions and kaons in matter arises from the term involving the vector current  $V_\mu$ , i.e.  $\text{Tr } \bar{B} i v_\mu \mathcal{D}^\mu B$ . In case of the pion, this term is identical to the well-known Weinberg-Tomozawa term [37]

$$\delta \mathcal{L}_{WT} = \frac{-1}{4f_\pi^2} (\bar{N} \vec{\tau} \gamma^\mu N) \cdot (\vec{\pi} \times (\partial_\mu \vec{\pi})) . \quad (10)$$

It contributes only to the isovector s-wave scattering amplitude and, therefore, does not contribute to the pion optical potential in isospin symmetric nuclear matter. This is *different* in case of the kaon, where we have

$$\delta \mathcal{L}_{WT} = \frac{-i}{8f^2} \left( 3(\bar{N} \gamma^\mu N)(K \overleftrightarrow{\partial}_\mu K) + (\bar{N} \vec{\tau} \gamma^\mu N)(K \vec{\tau} \overleftrightarrow{\partial}_\mu K) \right) . \quad (11)$$

The first term contributes to the isoscalar s-wave amplitude and, therefore, gives rise to an attractive or repulsive optical potential for  $K^-$  and  $K^+$  in symmetric nuclear matter. It turns out, however, that this leading order Lagrangian leads to an s-wave scattering length which is too repulsive as compared with experiment. Therefore, terms next to leading order in the chiral expansion are needed. Some of these involve the kaon-nucleon sigma term and thus are sensitive to the second difference between pions and kaons, namely the strength of the explicit symmetry breaking. The next to leading order effective kaon-nucleon Lagrangian can be written as [72]

$$\begin{aligned} \mathcal{L}_{\nu=2} = & \frac{\Sigma_{KN}}{f^2} (\bar{N} N) (\bar{K} K) + \frac{C}{f^2} (\bar{N} \vec{\tau} N) \cdot (\bar{K} \vec{\tau} K) \\ & + \frac{\tilde{D}}{f^2} (\bar{N} N) (\partial_t \bar{K} \partial_t K) + \frac{\tilde{D}'}{f^2} (\bar{N} \vec{\tau} N) \cdot (\partial_t \bar{K} \vec{\tau} \partial_t K) . \end{aligned} \quad (12)$$

The value of the kaon-nucleon sigma-term  $\Sigma_{KN} = \frac{1}{2}(m_q + m_s) \langle N | \bar{u}u + \bar{s}s | N \rangle$  depends on the strangeness content of the nucleon,  $y = 2 \langle N | \bar{s}s | N \rangle / \langle N | \bar{u}u + \bar{d}d | N \rangle \approx 0.1 - 0.2$ . Using the light quark mass ratio  $m_s/m \approx 29$ , one obtains  $370 < \Sigma_{KN} < 405$  MeV. The additional parameters,  $C, \tilde{D}, \tilde{D}'$  are then fixed by comparing with  $K^+$ -nucleon scattering data [72]. This so determined effective Lagrangian can then be used to predict the  $K^-$ -nucleon scattering amplitudes, and one obtains an attractive isoscalar s-wave scattering length in contradiction with experiments, where one finds a repulsive amplitude [74]. This discrepancy has been attributed to the existence of the  $\Lambda(1405)$  which is located below the  $K^-N$ -threshold. From the analysis of  $K^-p \rightarrow \Sigma\pi$  reactions it is known that this resonance couples strongly to the  $I = 0$   $K^-p$  state, and, therefore, leads to repulsion in the  $K^-p$  amplitude.

THE ROLE OF THE  $\Lambda(1405)$ 

Already in the sixties [75] there have been attempts to understand the  $\Lambda(1405)$  as a bound state of the proton and the  $K^-$ . In this picture, the underlying  $K^-$ -proton interaction is indeed attractive as predicted by the chiral Lagrangians but the scattering amplitude is repulsive only because a bound state is formed. This concept is familiar to the nuclear physicist from the deuteron, which is bound, because of the attractive interaction between proton and neutron. The existence of the deuteron then leads to a repulsive scattering length in spite of the attractive interaction between neutron and proton. Chiral perturbation theory is based on a systematic expansion of the S-matrix elements in powers of momenta and, therefore, effects which are due to the proximity of a resonance, such as the  $\Lambda(1405)$ , will only show up in terms of rather high order in the chiral counting. Thus, it is not too surprising that the first two orders of the chiral expansion predict the wrong sign of the  $K^-N$  amplitude. To circumvent this problem, the chiral perturbation calculation has been extended to either include an explicit  $\Lambda(1405)$  state [76] or to use the interaction obtained from the leading order chiral Lagrangian as a kernel for a Lippman-Schwinger type calculation, which is then solved to generate a bound state  $\Lambda(1405)$  [77].

This picture of the  $\Lambda(1405)$  as a  $K^-p$  bound state has recently received some considerable interest in the context of in-medium corrections. In ref. [78] it has been pointed out that in this picture as a result of the Pauli-blocking of the proton inside this bound state, the properties of the  $\Lambda(1405)$  would be significantly changed in the nuclear environment. With increasing density, its mass increases and the strength of the resonance is reduced (see Fig. 2). Because of this shifting and ‘disappearance’ of the  $\Lambda(1405)$  in matter, the  $K^-$  optical potential changes sign from repulsive to attractive at a density of about 1/4 of nuclear matter density (see Fig. 2) in agreement with a recent analysis of  $K^-$  atoms [79]. These findings have been confirmed in ref. [80]. This in-medium change of the  $\Lambda(1405)$  due to the Pauli blocking can only occur if a large fraction of its wave function is indeed that of a  $K^-$ -proton bound state. Of course one could probably allow for a small admixture of a genuine three quark state without changing the results for the measured kaon potentials. But certainly, if the  $\Lambda(1405)$  is mostly a genuine three quark state, the Pauli blocking should not affect its properties. Therefore, it would be very interesting to confirm the mass shift of the  $\Lambda(1405)$  for instance by a measurement of the missing mass spectrum of kaons in the reaction  $p + \gamma \rightarrow \Lambda(1405) + K^+$  at CEBAF. Thus, the

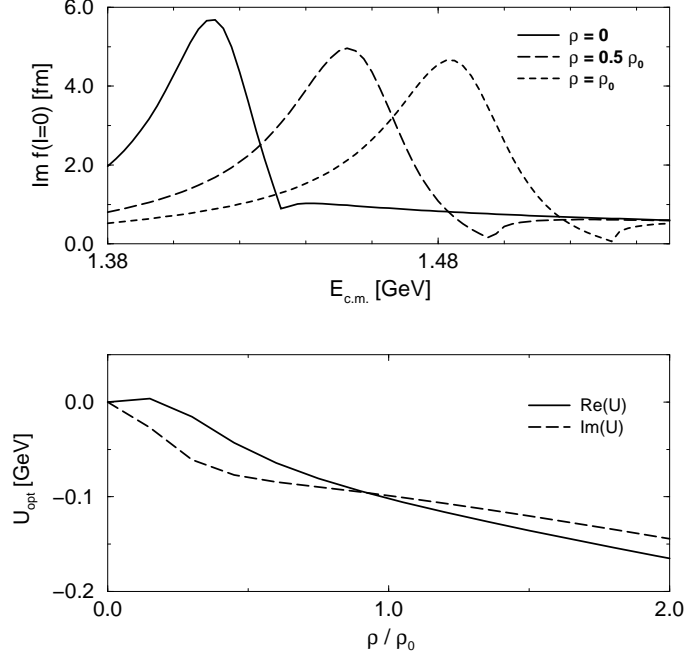


Figure 2: (a) Imaginary part of the  $I = 0$   $K^-$ -proton scattering amplitude for different densities. (b) real and Imaginary part of the  $K^-$  optical potential.

atomic nucleus provides a unique laboratory to investigate the properties of elementary particles. From our discussion it is clear that this in-medium change of the  $\Lambda(1405)$  is not related to the restoration of chiral symmetry.

#### EXPERIMENTAL RESULTS

Phenomenologically, the attractive optical potential for the  $K^-$  in nuclear matter is of particular interest because it can lead to a possible kaon condensation in neutron stars [72, 81]. This would limit the maximum mass of neutron stars to about one and a half solar masses and give rise to speculations about many small black holes in our galaxy [82]. Kaonic atoms, of course, only probe the very low density behavior of the kaon optical potential and, therefore, an extrapolation to the large



densities relevant for neutron stars is rather uncertain. Additional information about the kaon optical potential can be obtained from heavy ion collision experiments, where densities of more than twice nuclear matter density are reached.

Observables that are sensitive to the kaon mean-field potentials are the subthreshold production [83, 84, 85] as well as kaon flow [86, 87]. Given an attractive/repulsive mean-field potential for the kaons it is clear that the subthreshold production is enhanced/reduced. In case of the kaon flow an attractive interaction between kaons and nucleons aligns the kaon flow with that of the nucleons whereas a repulsion leads to an anti-alignment (anti-flow). However, both observables are also extremely sensitive to the overall reaction dynamics, in particular to the properties of the nuclear mean field and to reabsorption processes especially in case of the antikaons. Therefore, transport calculations are required in order to consistently incorporate all these effects.

In Fig. 3 we show the result obtained from such calculations for the  $K^+$  subthreshold production and flow together with experimental data (solid circles) from the KaoS collaboration [88] and from the FOPI collaboration [89] at GSI. Results are shown for calculations without any mean fields (dotted curves) as well as with a repulsive mean field obtained from the chiral Lagrangian (solid curves), which is consistent with a simple impulse approximation. The experimental data are nicely reproduced when the kaon mean-field potential is included. Although the subthreshold production of kaons can also be explained without any kaon mean field [90, 91], the assumption used in these calculations that a lambda particle has the same mean-field potential as a nucleon is not consistent with the phenomenology of hypernuclei [92]. Since it is undisputed that these observables are sensitive to the in-medium kaon potential, a systematic investigation including all observables should eventually reveal more accurately the strength of the kaon potential in dense matter. Additional evidence for a repulsive  $K^+$  potential comes from  $K^+$ -nucleus experiments. There the measured cross sections and angular distributions can be pretty well understood within a simple impulse approximation. Actually it seems that an additional (15%) repulsion is required to obtain an optimal fit to the data [93], which has been suggested as a possible evidence for a swelling of the nucleon size or a lowering of the omega meson mass in the nuclear medium [94].

As for the  $K^-$ , both chiral perturbation theory and dynamical models of the  $K^-$ -nucleon interaction [78] indicate the existence of an attractive mean field, so the same observables can be used [95, 96]. However, the measurement and interpretation of  $K^-$  observables is more difficult

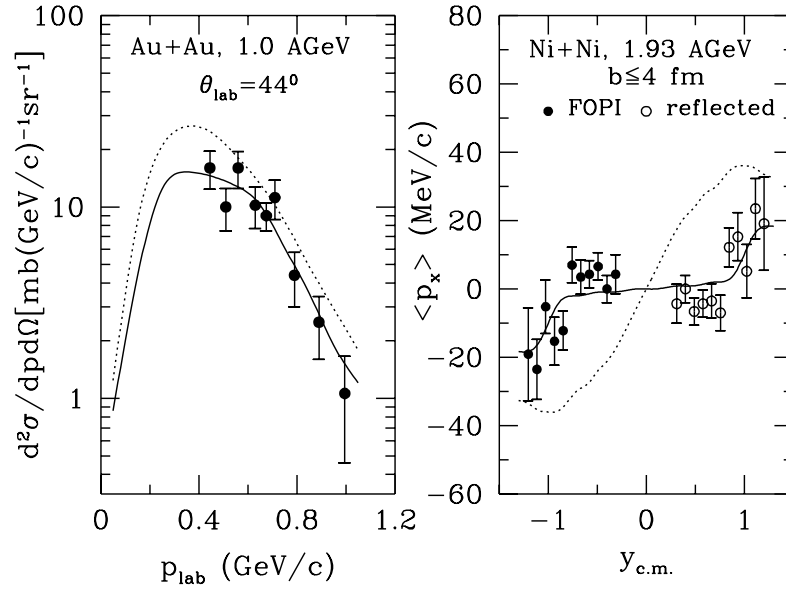


Figure 3: Kaon yield (left panel) and flow (right panel) in heavy ion collisions. The solid and dotted curves are results from transport model calculations with and without kaon mean-field potential, respectively. The data are from refs. [88, 89].

since it is produced less abundantly. Also reabsorption effects due to the reaction  $K^- N \rightarrow \Lambda \pi$  are strong, which further complicates the analysis. Nevertheless, the effect of the attractive mean field has been shown to be significant as illustrated in Fig. 4, where results from transport calculations with (solid curves) or without (dotted curves) attractive mean field for the antikaons [95] are shown. It is seen that the data (solid circles) on subthreshold  $K^-$  production [97] support the existence of an attractive antikaon mean-field potential. For the  $K^-$  flow, there only exist very preliminary data from the FOPI collaboration [98], which seem to show that the  $K^-$ 's have a positive flow rather an antifold, thus again consistent with an attractive antikaon mean-field potential.

Let us conclude this section on kaons by pointing out that the properties of kaons in matter are qualitatively described in chiral perturbation theory. Although some of the effects comes from the Weinberg-Tomozawa vector type interaction, which contributes because contrary to the pion the kaon has only one light quark, in order to explain the experimental data on kaon subthreshold production and flow in heavy ion collisions an additional attractive scalar interaction is required, which is related to the explicit breaking of chiral symmetry and higher order corrections in the chiral expansion.

### 2.3 *Etas*

The properties of etas are not only determined by consideration of chiral symmetry but also by the explicit breaking of the  $U_A(1)$  axial symmetry due to the axial anomaly in QCD. As a result, the singlet eta, which would be a Goldstone boson if  $U_A(1)$  were not explicitly broken, becomes heavy. Furthermore, because of SU(3) symmetry breaking – the strange quark mass is considerably heavier than that of up and down quark – the octet eta and singlet eta mix, leading to the observed particles  $\eta$  and  $\eta'$ . The mixing is such that the  $\eta'$  is mostly singlet and thus heavy, and the  $\eta$  is mostly octet and therefore has roughly the mass of the kaons. If, as has been speculated [99, 100], the  $U_A(1)$  symmetry is restored at high temperature due to the instanton effects, one would expect considerable reduction in the masses of the etas as well as in their mixing [101, 102]. A dropping eta meson in-medium mass is expected to provide a possible explanation for the observed enhancement of low transverse momentum etas in SIS heavy ion experiments at subthreshold energies [103]. However, an analysis of photon spectra from heavy ion collisions at SPS-energies puts an upper limit on the  $\eta/\pi^0$  ratio to be not more than 20 % larger for central than for peripheral collisions [104].

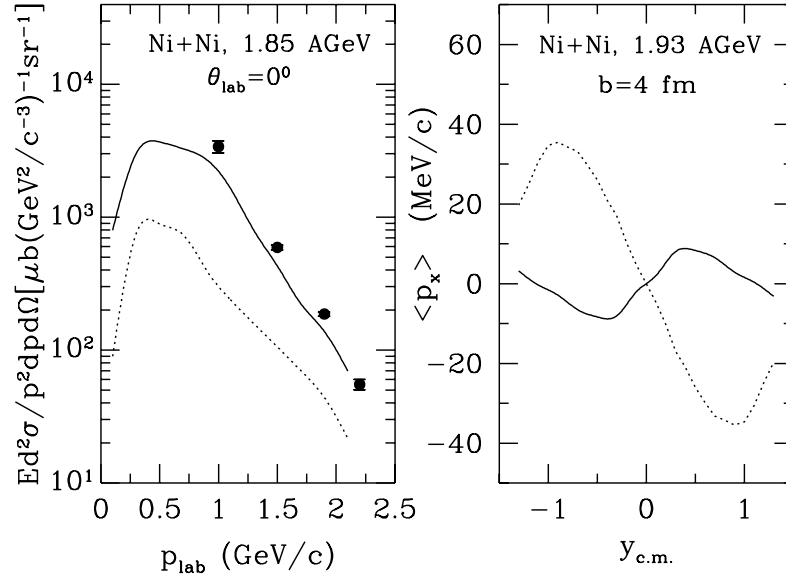


Figure 4: Antikaon yield (left panel) and flow (right panel) in heavy ion collisions. The solid and dotted curves are results from transport model calculations with and without kaon mean-field potential, respectively. The data are from ref. [97].

This imposes severe constraints on the changes of the  $\eta$  properties in hadronic matter.

We note that the restoration of chiral symmetry is important in the  $\eta - \eta'$  sector, because without the  $U_A(1)$  breaking both would be Goldstone bosons of an extended  $U(3) \times U(3)$  symmetry. However, it is still being debated what the effects precisely are.

### 3 VECTOR MESONS

Of all particles it is probably the  $\rho$ -meson which has received the most attention in regards of in-medium corrections. This is mainly due to the fact that the  $\rho$  is directly observable in the dilepton invariant mass spectrum. Also, since the  $\rho$  carries the quantum numbers of the conserved vector current, its properties are related to chiral symmetry and can, as we shall discuss, be investigated using effective chiral models as well as current algebra and QCD sum rules. That possible changes of the  $\rho$  can be observed in the dilepton spectrum measured in heavy ion collisions has been first demonstrated in ref. [105] and then studied in more detail in [106, 107]. The in-medium properties of the  $a_1$  are closely related to that of the  $\rho$  since they are chiral partners and their mass difference in vacuum is due to the spontaneous breaking of chiral symmetry [108]. Unfortunately, there is no direct method to measure the changes of the  $a_1$  in hadronic matter. The  $\omega$ -meson on the other hand is a chiral singlet, and the relation of its properties to chiral symmetry is thus not so direct. However, calculations based on QCD sum rules predict also changes in the  $\omega$  mass as one approaches chiral restoration. Observationally, the  $\omega$  can probably be studied best due to its rather small width and its decay channel into dileptons. Finally, there is the  $\phi$  meson. In an extended SU(3) chiral symmetry, the  $\phi$  and the  $\omega$  are a superposition of the singlet and octet states with nearly perfect mixing, i.e. the  $\phi$  contains only strange quarks whereas the  $\omega$  is made only out of light quarks. Both QCD sum rules [109] and effective chiral models [110] predict a lowering of the  $\phi$ -meson mass in medium.

The question on whether or not masses of light hadrons change in the medium has received considerable interest as a result of the conjecture of Brown and Rho [8], which asserts that the masses of all mesons, with the exception of the Goldstone Bosons, should scale with the quark condensate. While the detailed theoretical foundations of this conjecture are still being worked on [111, 112, 113] the basic argument of Brown and Rho is as follows. Hadron masses, such as that of the  $\rho$ -meson,

violate scale invariance, which is a symmetry of the classical QCD Lagrangian. In QCD scale invariance is broken on the quantum level by the so-called trace anomaly (see e.g. [114]), which is proportional to the Gluon condensate. Thus one could imagine that with the disappearance of the gluon condensate, i.e. the bag pressure, scale invariance is restored, which on the hadronic level implies that hadron masses have to vanish. Therefore, one could argue that hadron masses should scale with the bag-pressure as originally proposed by Pisarski [115]. But the conjecture of Brown and Rho goes even further. They assume that the gluon condensate can be separated into a hard and soft part, the latter of which scales with the quark condensate and is also responsible for the masses of the light hadrons. This picture finds some support from lattice QCD calculations in that the gluons condensate drops by about 50% at the chiral phase transition [116]. To what extent this is also reflected in changes in hadron masses is not clear at the moment (see section 5), although one should mention that the rise in the entropy density close to the critical temperature can be explained if one assumes the hadron masses to scale with the quark condensate [116]. Another aspect of the Brown and Rho scaling is that once the scaling hadron masses are introduced in the chiral Lagrangian, only tree-level diagrams are needed as the contribution from higher order diagrams is expected to be suppressed. In their picture, Goldstone bosons are not subject to this scaling, since they receive their mass from the explicit chiral symmetry breaking due to finite current quark masses, which are presumably generated at a much higher (Higgs) scale.

### 3.1 *The rho meson*

Since the  $\rho$  is a vector meson, it couples directly to the isovector current which then results in the direct decay of the  $\rho$  into virtual photons, i.e. dileptons. Consequently, properties of the  $\rho$  meson can be investigated by studying two-point correlation functions of the isovector currents, i.e.,

$$\Pi_{\mu\nu}(q) = i \int e^{iqx} \langle T J_\mu(x) J_\nu(0) \rangle_\rho d^4x. \quad (13)$$

The masses of the rho meson and its excitations ( $\rho' \dots$ ) correspond to the positions of the poles of this correlation function. This is best seen if one assumes that the current field identity [117] holds, namely that the current operator is proportional to the  $\rho$ -meson field. In this case the above correlation function is identical, up to a constant, to the

$\rho$ -meson propagator. The imaginary part of this correlation function is also directly proportional to the electron-positron annihilation cross section [118], where the  $\rho$ -meson is nicely seen. In addition, at higher center-of-mass energies, one sees a continuum in the electron-positron annihilation cross section, which corresponds to the excited states of the rho mesons as well as to the onset of perturbative QCD processes.

In-medium changes of the  $\rho$  meson can be addressed theoretically by evaluating the current-current correlator in the hadronic environment. The current-current correlator can be evaluated either in effective chiral models, or using current algebra arguments, or directly in QCD. In the latter, one evaluates the correlator in the deeply Euclidean region ( $q^2 \rightarrow -\infty$ ) using the Wilson expansion, where all the long distance physics is expressed in terms of vacuum expectation values of quark and gluon operators, the so-called condensates. Dispersion relations are then used to relate the correlator in the Euclidean region to that for positive  $q^2$ , where the hadronic eigenstates are located. One then assumes a certain shape for the phenomenological spectral functions, typically a delta function, which represents the bound state, and a continuum, which represents the perturbative regime. These so called QCD sum rules, therefore, relate the observable hadronic spectrum with the QCD vacuum condensates (for a review of the QCD sum-rule techniques see e.g. [119]). These relations can then either be used to determine the condensates from measured hadronic spectra, or, to make predictions about changes of the hadronic spectrum due to in-medium changes of the condensates.

Similarly, one can study the properties of  $a_1$  meson in the nuclear medium through the axial vector correlation function. Then, once chiral symmetry is restored, there should be no observable difference between left-handed and right-handed or equivalently vector and axial vector currents. Consequently, the vector and axial vector correlators should be identical. Often, this identity of the correlators is identified with the degeneracy of the  $\rho$  and  $a_1$  mesons in a chirally symmetric world. This, however, is not the only possibility, as was pointed out by Kapusta and Shuryak [120]. There are at least three qualitatively different scenarios, for which the vector and axial vector correlator are identical (see Fig. 5).

1. The masses of  $\rho$  and  $a_1$  are the same. The value of the common mass, however, does not follow from chiral symmetry arguments alone.
2. The mixing of the spectral functions, i.e, both the vector and

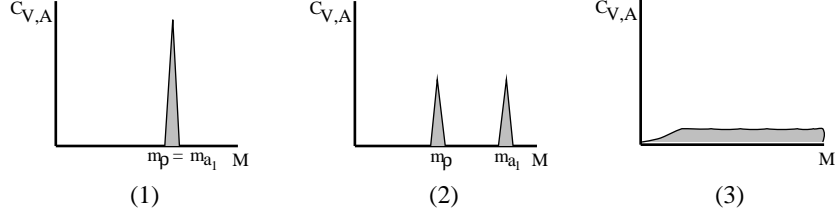


Figure 5: Several possibilities for the vector and axial-vector spectral functions in the chirally restored phase.

axial-vector spectral functions have peaks of similar strength at both the mass of the  $\rho$  and the mass of the  $a_1$ .

3. Both spectral functions could be smeared over the entire mass range. Because of thermal broadening of the mesons and the onset of deconfinement, the structure of the spectral function may be washed out, and it becomes meaningless to talk about mesonic states.

#### FINITE TEMPERATURE

At low temperatures, where the heat bath consists of pions only, one can employ current algebra as well as PCAC to obtain an essentially model-independent result for the properties of the  $\rho$ . Using this technique Dey, Eletsky and Ioffe [121] could show that to leading order in the temperature,  $T^2$ , the mass of the  $\rho$ -meson does not change. Instead one finds an admixture of the axial-vector correlator, i.e. that governed by the  $a_1$ -meson. Specifically to this order the vector correlator is given by

$$C_V(q, T) = (1 - \epsilon) C_V(q, T = 0) + \epsilon C_A(q, T = 0) + \mathcal{O}(T^4), \quad (14)$$

with  $\epsilon = T^2/(6f_\pi^2)$ . Here  $C_V, C_A$  stand for the vector-isovector and axial-vector-isovector correlation functions, respectively. One should note, that to the same order the chiral condensate is reduced [59],

$$\frac{\langle \bar{q}q \rangle_T}{\langle \bar{q}q \rangle_0} = 1 - \frac{T^2}{8f_\pi^2} + \mathcal{O}(T^4). \quad (15)$$

Therefore, to leading order a drop in the chiral condensate does not affect the mass of the  $\rho$ -meson but rather reduces its coupling to the vector current and in particular induces an admixture of the  $a_1$ -meson.



This finding is at variance with the Brown-Rho scaling hypothesis. The reason for this difference is not yet understood. The admixture of the axial correlator is directly related with the onset of chiral restoration. If chiral symmetry is restored, the vector and axial-vector correlators should be identical. The result of Dey et al. suggests that this is achieved by a mixing of the two instead of a degeneracy of the  $\rho$  and  $a_1$  masses. If the mixing is complete, i.e.  $\epsilon = 1/2$ , then the extrapolation of the low temperature result (14) would give a critical temperature of  $T_c = \sqrt{3}f_\pi \simeq 164 \text{ MeV}$ , which is surprisingly close to the value given by recent lattice calculations.

Corrections to the order  $T^4$  involve physics beyond chiral symmetry. As nicely discussed in ref. [122], to this order the contributions can be separated into two distinct contributions. The first arises if a pion from the heat bath couples via a derivative coupling to the current under consideration. In this case the contribution is proportional to the invariant pion density

$$n_\pi = \int \frac{d^3q}{2\omega(2\pi)^3} e^{-\omega/T} \sim T^2, \quad (16)$$

and the square of the pion momentum  $q^2 \sim T^2$ . A typical example would be for instance the self-energy correction to the  $\rho$  meson from the standard two-pion loop diagram involving the p-wave  $\pi\pi\rho$  coupling. The other contribution comes from interactions of two pions from the heat bath with the current, without derivative couplings. This is proportional to the square of the invariant pion density and thus  $\sim T^4$  as well. These pure density contributions again can be evaluated in a model independent fashion using current algebra techniques and, as before, do not change the mass of the  $\rho$ -meson, but change the coupling to the current and induce a mixing with the axial-vector correlator. Actually, it can be shown [122] that to all orders in the pion density these pure density effects do not change the  $\rho$  mass but induce mixing and reduce the coupling. At the same time, however, these contributions reduce the chiral condensate.

The contributions due to the finite pion momenta have been estimated in [122] to give a downward shift of both the  $\rho$  and  $a_1$  mass of about

$$\frac{\delta m}{m} \simeq 10\%. \quad (17)$$

for temperatures of  $\sim 150 - 200 \text{ MeV}$ . We should point out, however, that in this analysis those changes in the masses arise from Lorentz-nonscalar condensates in the operator product expansion. Therefore,

they are not directly related to the change of the Lorentz-scalar chiral-condensate.

Effective chiral models also have been used to explore the order  $T^4$  corrections. Ref. [123] reports that the mass of the  $\rho$  drops whereas that of the  $a_1$  increases to this order. This is somewhat at variance with the findings of [124] where no drop in the mass of the  $\rho$  but a decrease in the  $a_1$  mass has been found. In this calculation, however, the terms leading to the order  $T^4$  changes have not been explicitly identified but rather a calculation to one loop order has been carried out. Presently, the difference in these results is not understood. Also the difference between the analysis of [122] and the effective chiral models is not resolved yet. The analysis of [122] uses dispersion relations to relate phenomenological space-like photon-pion amplitudes with the time-like ones needed to calculate the mass shift. The effective Lagrangian methods, on the other hand, rely on pion-pion scattering data as well as measured decay width in order to fix their model parameters. One would think that both methods should give a reasonably handle on the leading order momentum-dependent couplings.

At higher temperatures as well as at finite density vacuum effects due to the virtual pair correction or the nucleon-antinucleon polarization could become important and they tend to reduce vector meson masses [125]. Because of the large mass of the nucleon, these corrections however, do not affect the leading temperature result  $\sim T^2$  discussed previously. Also this approach assumes that the physical vacuum consists of nucleon-antinucleon rather than quark-antiquark fluctuations. Whether this is the correct picture is, however, not yet resolved.

#### FINITE DENSITY

Since the density effect on the chiral condensate is much stronger than that of the temperature, one expects the same for the rho meson mass. But the situation is more complicated at finite density as one cannot make use of current algebra arguments and thus model-independent result as the one discussed previously are not available at this time. Present model calculations, however, disagree even on the sign of the mass shift. One class of models [126, 127, 128, 129, 130] considers the  $\rho$  as a pion-pion resonance and calculates its in-medium modifications due to those for the pions as discussed in section 2.1. These calculations typically show an increased width of the rho since it now can also decay into pion-nucleon-hole or pion-delta-hole states. At the same time this leads to an increased strength below the rho meson mass. The mass

of the rho meson, defined as the position where the real part of the correlation function goes through zero, is shifted only very little. Most calculations give an upward shift but also a small downward shift has been reported [130, 131]. This difference seems to depend on the specific choice of the cutoff functions for the vertices involved [132], and thus is model dependent. However, the imaginary part of the correlation function, the relevant quantity for the dilepton measurements, hardly depends on a small upward or downward shift of the rho. The important and apparently model independent feature is the increased strength at low invariant masses due to the additional decay-channels available in nuclear matter.

Calculation using QCD-sum rules [133] predict a rather strong decrease of the  $\rho$ -mass with density ( $\sim 20\%$  at nuclear matter density), which is similar to the much discussed Brown-Rho scaling [8]. Here, as in the finite temperature case, the driving term is the four quark condensate which is assumed to factorize

$$\langle (\bar{q}\gamma_\mu\lambda^a q)(\bar{q}\gamma^\mu\lambda^a q) \rangle_\rho \approx -\frac{16}{9}\langle\bar{q}q\rangle_\rho^2. \quad (18)$$

To which extent this factorization is correct at finite density is not clear. Also, when it comes to the parameterization of the phenomenological spectral distribution, these calculations usually assume the standard pole plus continuum form with the addition of a so called Landau damping contribution at  $q^2 = 0$  [133]. The additional strength below the mass of the rho as predicted by the previously discussed models is usually ignored. An attempt, however, has been made to see if a spectral distribution obtained from the effective models described above does saturate the QCD sum rules [134]. These authors found that they could only saturate the sum rule if they assumed an additional mass shift of the rho-meson peak downwards by  $\sim 140$  MeV. However, in this calculation only leading order density corrections to the quark condensates whereas infinity order density effects have been taken into account in order to calculate the spectrum. Another comparison with the QCD sum rules has recently been carried out in ref. [131] where a reasonable saturation of the sum-rule is reported using a similar model for the in medium correlation function.

There are also attempts to use photon-nucleon data in order to estimate possible mass shifts of the  $\rho$  and  $\omega$  mesons. In ref. [135] a simple pion and sigma exchange model is used in order to fit data for photoproduction of  $\rho$ - and  $\omega$ -mesons. Assuming vector dominance, this model is then used to calculate the self-energy of these vector mesons in nuclear

matter. The authors find a downward shift of the  $\rho$  of about 18% at nuclear matter density, in rough agreement with the prediction from QCD sum rules. Quite to the contrary, ref. [136] using photoabsorption data and dispersion relations find an upwards shift of 10 MeV or 50 MeV for the longitudinal and transverse part of the  $\rho$ , respectively. This result, however, is derived for a  $\rho$ -meson which is not at rest in the nuclear matter frame, and, therefore a direct comparison of the two predictions is not possible.

Very recently Friman and Pirner have pointed out that the  $\rho$ -meson couples very strongly with the  $N^*(1720)$  resonance [31, 137]. The coupling is of p-wave nature and, similarly to the pion coupling to delta-hole states, the  $\rho$  may mix with  $N^*$ -nucleon-hole states resulting in a modified dispersion relation for the  $\rho$ -meson in nuclear matter. Since the coupling is p-wave, only  $\rho$ -mesons with finite momentum are modified and shifted to lower masses. This momentum dependence of the low mass enhancement in the dilepton spectrum is an unique prediction which can be tested in experiment.

#### EXPERIMENTAL RESULTS

First measurement of dileptons in heavy ion collisions have been carried out by the DLS collaboration [53, 54] at the BEVALAC. The first published data based on a limited data set could be well reproduced using state of the art transport models [106, 138, 139] without any additional in-medium corrections. However, a recent reanalysis [140] including the full data set seems to show a considerable increase over the originally published data. It remains to be seen if these new data can also be understood without any in-medium corrections to hadronic properties. Recently, low mass dilepton data taken at the CERN SPS have been published by the CERES collaboration [142]. Their measurements for p+Be and p+Au could be well understood within a hadronic cocktail, which takes into account the measured particle yields from p+p experiments and their decay channels into dileptons. In case of the heavy ion collision, S+Pb, however, the hadronic cocktail considerably under-predicted the measured data in particular in the invariant mass region of  $300 \text{ MeV} \leq M_{\text{inv}} \leq 500 \text{ MeV}$ . A similar enhancement has also been reported by the HELIOS-3 collaboration [143].

Of course it is well-known that in a heavy ion collisions at SPS-energies a hadronic fireball consisting predominantly of pions is created. These pions can pairwise annihilate into dileptons giving rise to an additional source, which has not been included into the hadronic cocktail.

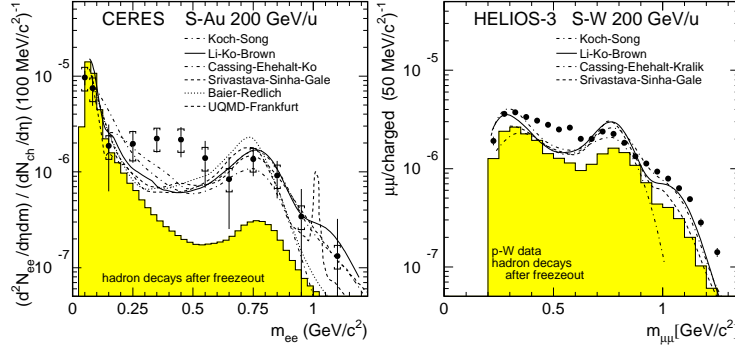


Figure 6: Dilepton invariant mass spectrum from S+Au and S+W collisions at 200 GeV/nucleon without dropping vector meson masses. The experimental statistical errors are shown as bars and the systematic errors are marked independently by brackets. (The figure is from ref. [141])

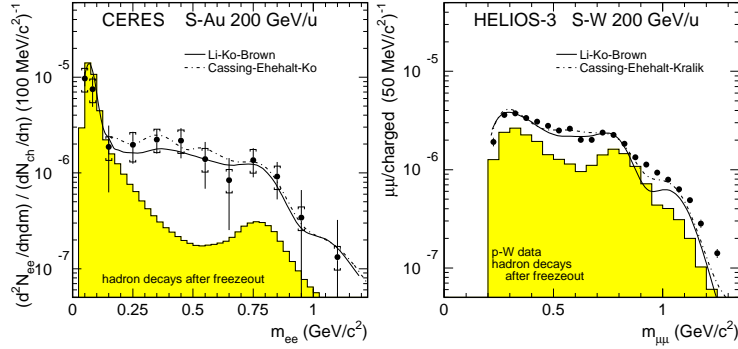


Figure 7: Same as Fig. 6 with dropping vector meson masses. (The figure is from ref. [141])

While the pion annihilation contributes to the desired mass range, many detailed calculations [65, 144, 145], which have taken this channel into account, still underestimate the data by about a factor of three as shown Fig. 6. Also, additional hadronic processes, such as  $\pi\rho \rightarrow \pi e^+ e^-$ , have been taken into account [146, 147]. However, ‘conservative’ calculations could at best reach the lower end of the sum of statistical and systematic error bars of the CERES data. Furthermore, in-medium corrections to the pion annihilation process together with corrections to the  $\rho$ -meson spectral distribution have been considered using effective chiral models [64, 148]. But, these in-medium corrections, while enhancing the contribution of the pion annihilation channel somewhat, are too small to reproduce the central data points. Only models which allow for a dropping of the  $\rho$ -meson mass give enough yield in the low mass region [144, 149, 150, 151] as shown in Fig. 7. In these models, the change of the rho meson mass is obtained from either the Brown-Rho scaling [149, 150, 151] or the QCD sum rules [144]. Furthermore, the vector dominance is assumed to be suppressed in the sense that the pion electromagnetic form factor, which in free space is dominated by the rho meson and is proportional to the square of the rho meson mass, is reduced as a result of the dropping rho meson mass. At finite temperature, such a suppression has been shown to exist in both the hidden gauge theory [152] and the perturbative QCD [153], where the pion electromagnetic form factor is found to decrease at the order  $T^2$ . This effect is related to the mixing between the vector-isovector and axial-vector-isovector correlators at finite temperature we discussed earlier. In most studies of dilepton production without a dropping rho meson mass, this effect has not been included. Although the temperature reached in heavy ion collisions at SPS energies is high and the ratio of pions to baryons in the final state is about 5 to 1, the authors in [144, 149, 150, 151] found that they had to rely on the nuclear density rather than on the temperature in order to reproduce the data.

There are also preliminary data from the Pb+Au collisions at 150 GeV/nucleon [141], which also seem to be consistent with the dropping in-medium rho meson scenario as well. Unfortunately, the experimental errors in this case are even larger than in the S+Au data and thus these new, preliminary data do not further discriminate between the dropping rho-mass and more conventional scenarios. Thus, additional data with reduced error bars are needed before firm conclusions about in-medium changes of the rho meson mass can be drawn. Although the data from the HELIOS-3 collaboration shown in the right panels of Figs. 6 and 7 do have very small errors, these are only the statistical ones while the

systematic errors are not known.

To summarize this section on the rho meson, our present theoretical understanding for possible in-medium changes at low temperatures and vanishing baryon density indicates that one expects no (to leading order) and a small (to next to leading order) changes of the  $\rho$ -mass. At the same time to leading order in the temperature, the chiral condensate is reduced. Therefore, a direct connection between changes in the chiral condensate and the mass of the  $\rho$  meson has not been established as assumed in the Brown-Rho scaling. However, it may well be that the latter is only valid at higher temperature near the chiral phase transition. At finite density additional effects such as nuclear many-body excitations come into play, which make the situation much more complicated. Nevertheless, the observed enhancement of low mass dileptons from CERN-SPS heavy ion collisions can best be described by a dropping rho meson mass in dense matter. Finite density and zero temperature is another area where possible changes of the vector mesons can be measured in experiment. Photon, proton and pion induced dilepton production from nuclei will soon be measured at CEBAF [154] and at GSI [155]. In principle these measurements should be able to determine the entire spectral distribution. If the predictions of the Brown-Rho scaling and the QCD sum rules are correct, mass shifts of the order of 100 MeV should occur, which would be visible in these experiments. Furthermore, by choosing appropriate kinematics, the properties of the  $\rho$  meson at rest as well as at finite momentum with respect to the nuclear rest frame can be measured.

### 3.2 *The omega meson*

The omega meson couples to the isoscalar part of the electromagnetic current. In QCD sum rules it is also dominated by the four quark condensate. If the latter is reduced in medium, then the omega meson mass also decreases. Assuming factorization for the four quark condensate, Hatsuda and Lee [133] showed that the change of the omega meson mass in dense nuclear matter would be similar to that of the rho meson mass. The dropping omega meson mass in medium can also be obtained by considering the nucleon-antinucleon polarization in medium [125, 156, 157] as in the case of the rho meson. On the other hand, using effective chiral Lagrangians, one finds at finite temperature an even smaller change in the mass of the omega as compared to that of the rho [124]. Because of its small decay width, which is about an order of magnitude smaller than that of rho meson, most dileptons from omega

decay in heavy ion collisions are emitted after freeze out, where the medium effects are negligible. However, with appropriate kinematics an omega can be produced at rest in nuclei in reactions induced by the photon, proton and pion. By measuring the decay of the omega into dileptons allows one to determine its properties in the nuclear matter. Such experiments will be carried out at CEBAF [154] and GSI [155].

### 3.3 *The phi meson*

For the phi meson, the situation is less ambiguous than the rho meson as both effective chiral Lagrangian and QCD sum-rule studies predict that its mass decreases in medium. Based on the hidden gauge theory, Song [110] finds that at a temperature of  $T = 200$  MeV the phi meson mass is reduced by about 20 MeV. The main contribution is from the thermal kaon loop. In QCD sum rules, the reduction is much more appreciable, i.e., about 200 MeV at the same temperature [109]. The latter is due to the significant decrease of the strange quark condensate at finite temperature as a result of the abundant strange hadrons in hot matter. Because of the relative large strange quark mass compared to the up and down quark masses, the phi meson mass in QCD sum rules is mainly determined by the strange quark condensate instead of the four quark condensate as in the case of rho meson mass. However, the temperature dependence of the strange quark condensate in [109] is determined from a non-interacting gas model, so effects due to interactions, which may be important at high temperatures, are not included. In [110], only the lowest kaon loop has been included, so the change of kaon properties in medium is neglected. As shown in [158], this would reduce the phi meson mass if the kaon mass becomes small in medium. Also, vacuum effects in medium due to lambda-antilambda polarization has also been shown to reduce the phi meson in-medium mass [159].

QCD sum rules have also been used in studying phi meson mass at finite density [133], and it is found to decrease by about 25 MeV at normal nuclear matter density.

Current experimental data on phi meson mass from measuring the  $K\bar{K}$  invariance mass spectra in heavy ion collisions at AGS energies do not show a change of the phi meson mass [160]. This is not surprising as these kaon-antikaon pairs are from the decay of phi mesons at freeze out when their properties are the same as in free space. If a phi meson decays in medium, the resulting kaon and antikaon would interact with nucleons, so their invariant mass is modified and can no longer be used to reconstruct the phi meson. However, future experiments on measuring



dileptons from photon-nucleus reactions at CEBAF [154] and heavy ion collisions at GSI [155] will provide useful information on the phi meson properties in dense nuclear matter. For heavy ion collisions at RHIC energies, matter with low-baryon chemical potential is expected to be formed in central rapidities. Based on the QCD-sum rule prediction for the mass shift of the phi meson it has been suggested that a low mass phi peak at  $\sim 880$  MeV besides the normal one at 1.02 GeV appears in the dilepton spectrum if a first-order phase transition or a slow cross-over between the quark-gluon plasma and the hadronic matter occurs in the collisions [161, 162]. The low-mass phi peak is due to the nonnegligible duration time for the system to stay near the transition temperature compared with the lifetime of a phi meson in vacuum, so the contribution to dileptons from phi meson decays in the mixed phase becomes comparable to that from their decays at freeze out. Without the formation of the quark-gluon plasma, the low-mass phi peak is reduced to a shoulder in the dilepton spectrum. Thus, one can use this double phi peaks in the dilepton spectrum as a signature for identifying the quark-gluon plasma to hadronic matter phase transition in ultrarelativistic heavy ion collisions.

## 4 BARYONS

As far as the in medium properties of the baryons are concerned most is known about the nucleon. But also the properties of the hyperons such as the  $\Lambda$  can be determined in the medium by studying hypernuclei. Also the  $\Lambda(1405)$ , which we have discussed in connection with the  $K^-$  optical potential, is another example for in-medium effects of baryons. In the following we will limit ourself on a brief discussion of how the nucleon properties can be viewed in the context of chiral symmetry.

As mentioned briefly in the introduction, medium effects on a nucleon due to many-body interactions have long been studied, leading to an effective mass, which is generally reduced as a result of the finite range of the nucleon interaction and higher order effects. Its relation to chiral symmetry restoration is best seen through the QCD sum rules [2, 3, 163, 164, 165, 166, 167]. Using in-medium condensates, it has been shown that the change of scalar quark condensate in medium leads to an attractive scalar potential which reduces the nucleon mass, while the change of vector quark condensate leads to a repulsive vector potential which shifts its energy. The nucleon scalar and vector self-energies in

this study are given, respectively, by

$$\begin{aligned}\Sigma_S &\approx -\frac{8\pi^2}{M_B^2}(\langle\bar{q}q\rangle_\rho - \langle\bar{q}q\rangle_0) \approx -\frac{8\pi^2}{M_B^2}\frac{\Sigma_{\pi N}}{m_u + m_d}\rho_N, \\ \Sigma_V &\approx \frac{64\pi^2}{3M_B^2}\langle q^+q\rangle_\rho = \frac{32\pi^2}{M_B^2}\rho_N,\end{aligned}\tag{19}$$

where the Borel mass  $M_B$  is an arbitrary parameter. With  $M_B \approx m_N$ , and  $m_u + m_d \approx 11$  MeV, these self-energies have magnitude of a few hundred MeV at normal nuclear density. These values are similar to those determined from both the Walecka model [168] and the Dirac-Brueckner-Hartree-Fock (DBHF) approach [169] based on the meson-exchange nucleon-nucleon interaction. Experimental evidences for these strong scalar and vector potentials have been inferred from the proton-nucleus scattering at intermediate energies [170] via the Dirac phenomenology [171, 172, 173, 174] in which the Dirac equation with scalar and vector potentials is solved. The importance of medium effects on the nucleon mass can be seen from the significant decrease of the Q-value for the reaction  $NN \rightarrow NNN\bar{N}$  in nuclear medium as a result of the attractive scalar potential. This effect has been included in a number of studies based on transport models. In Fig. 8, theoretical results from these calculations for the antiproton differential cross section in Ni+Ni collisions at 1.85 GeV/nucleon [175, 176, 177] and C+Cu collisions at 3.65 GeV/nucleon [178] are compared with the experimental data from GSI [179] and Dubna [180]. In [175, 176, 177], the relativistic transport model has been used with the antiproton mean-field potential obtained from the Walecka model. The latter has a value in the range of -150 to -250 MeV at normal nuclear matter density. The results of ref. [178] are based on the nonrelativistic Quantum Molecular Dynamics with the produced nucleon and antinucleon masses taken from the Nambu–Jona-Lasinio model. Within this framework, these studies thus show that in order to describe the antiproton data from heavy-ion collisions at subthreshold energies it is necessary to include the reduction of both nucleon and antinucleon masses in nuclear medium. Even at AGS energies, which are above the antiproton production threshold in NN interaction, the medium effects on antiproton may still be important [181]. Indeed, a recent study using the Relativistic Quantum Molecular Dynamics shows that medium modifications of the antiproton properties are important for a quantitative description of the experimental data [182]. We note, however, a better understanding of antiproton annihilation is needed, as only about 5-10% of produced antiprotons can

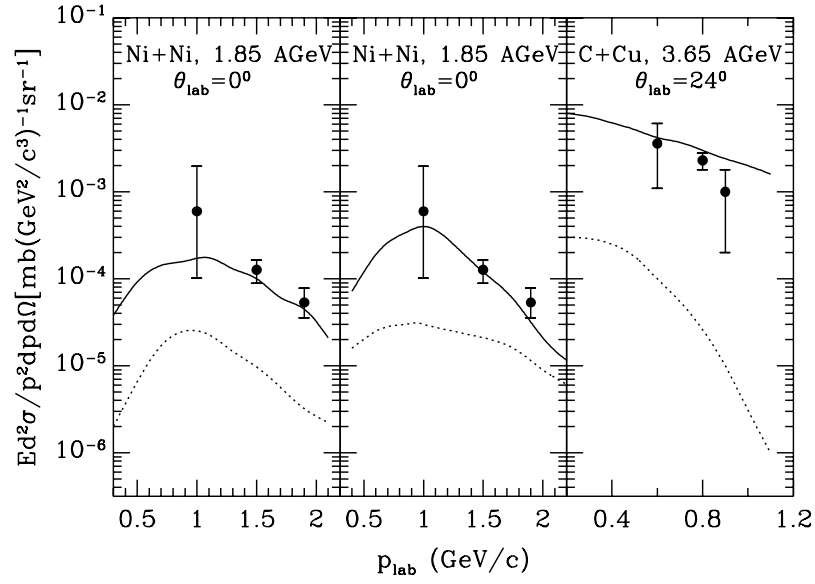


Figure 8: Antiproton momentum spectra from Ni+Ni collisions at 1.85 GeV/nucleon, and C+Cu collisions at 3.85 GeV/nucleon. The left, middle, and right panels are from Refs. [175], [177], and [178], respectively. The solid and dotted curves are from transport model calculations with free and dropping antiproton mass, respectively. The experimental data from ref. [179] for Ni+Ni collisions and from ref. [180] for C+Cu collisions are shown by solid circles.

survive, in order to determine more precisely the antiproton in-medium mass from heavy ion collisions.

## 5 RESULTS FROM LATTICE QCD CALCULATIONS

Another source of information about possible in medium changes are lattice QCD-calculations. For a review see e.g. [1]. Presently lattice calculations can only explore systems at finite temperature but vanishing baryon chemical potential. Therefore, their results are not directly applicable to present heavy ion experiments, where system at finite baryon chemical potential are created. But future experiments at RHIC and LHC might succeed in generating a baryon free region. But aside from the experimental aspects, lattice results provide an important additional source where our model understanding can be tested. Lattice calculation are usually carried out in Euclidean space, i.e. in a space with imaginary time. As a consequence plane waves in Minkowski-space translate into decaying exponentials in Euclidean space. At zero temperature, masses of the low lying hadrons are extracted from two point functions, which carry the quantum numbers of the hadron under consideration.

$$C(\tau) = \int d^3x \langle J(x, \tau) J(0) \rangle \sim \sum_i \alpha_i^2 \cdot \exp(-m_i \tau) \quad (20)$$

Here, the sum goes over all hadronic states which carry the quantum numbers of the operator  $J$ . At large imaginary times  $\tau$  only the state with the lowest mass survives, and its mass can be determined from the exponential slope. Zero temperature lattice calculation by now reproduce the hadronic spectrum to a remarkable accuracy [183]. If one wants to extract masses at finite temperature, however, things become more complicated. The reason is that finite temperature on the lattice is equivalent to requiring periodic or anti-periodic boundary conditions in the imaginary time direction for bosons or fermions, respectively. Therefore, one cannot study the correlation functions at arbitrary large  $\tau$  but is restricted to  $\tau_{max} = 1/T$ , where  $T$  is the temperature under consideration. As a result, the lowest lying states cannot be projected out so easily. This has led people to study the correlation functions as a function of the spatial distance, where no restriction to the spatial extent exist, to extract so-called screening masses. At zero temperature

the Euclidean time and space direction are equivalent and the screening masses are identical to the actual hadron masses. By analyzing the correlation function in the spatial direction one essentially measures the range of a virtual particle emission. As already demonstrated by Yukawa, who deduced the mass of the pion from the range of the nuclear interaction, this can be used to extract particle masses. Screening masses, however, are not very useful if one wants to study hadronic masses at high temperature, close to  $T_c$ . At these temperatures, the spatial correlators are dominated by the trivial contribution from free independent quarks, and the resulting screening masses simply turn out to be [184]

$$M_{\text{screen}} = n\pi T, \quad (21)$$

where  $n$  is the number of quarks of a given hadron, i.e.  $n = 2$  for mesons and  $n = 3$  for baryons. Only for the pion and sigma meson, lattice calculation found a significant deviation from this simple behavior, which is due to strong residual interactions in these channel. All other hadrons considered so far, such as the  $\rho$  and  $a_1$  as well as the nucleon, exhibit the above screening mass [185, 186]. These trivial contributions from independent quarks are absent, however, in the time-like correlator. There has been one attempt to extract meson masses from the time-like correlators in four flavor QCD [187]. In this case it has been found that the mass of the  $\rho$ -meson does not change significantly below  $T_c$ . Above the critical temperature, on the other hand, the correlation function appeared to be consistent with one of noninteracting quarks, indicating the onset of deconfinement.

There have been attempts to extract the mass of the scalar  $\sigma$  meson as well as that of the pion from so called susceptibilities [188]. These are nothing else than the total four-volume integral over the two-point functions. If the two-point function is dominated by one hadronic pole, then this integral should be inversely proportional to the square of the mass of the lightest hadron under consideration. Using this method, the degeneracy of the pion and sigma mass close to  $T_c$  has been demonstrated.

Finally, let us note that a scenario where all hadron masses scale linearly with the chiral condensate is consistent with strong rise in the energy and entropy density around the phase transition as observed in Lattice calculations [116].

## 6 SUMMARY

In this review, hadron properties, particularly their masses, in the nuclear medium have been discussed. Both conventional many-body interactions and genuine vacuum effects due to chiral symmetry restoration have effects on the hadron in-medium properties.

For Goldstone bosons, which are directly linked to the spontaneously broken chiral symmetry, we have discussed the pion, kaons, and etas. For the pion, its mass is only slightly shifted in the medium due to the small s-wave interactions as a result of chiral symmetry. On the other hand, the strong attractive p-wave interactions due to the  $\Delta(1230)$  and the  $\rho(770)$  lead to a softening of the pion dispersion relation in medium. These effects result from many-body interactions rather from chiral symmetry. Phenomenologically, the small change in the pion in-medium mass does not seem to have any observable effects. However, the effects due to the softened pion in-medium dispersion relation may be detectable through the enhanced low transverse momentum pions in heavy ion collisions.

For kaons, medium effects due to the Weinberg-Tomozawa vector type interaction, which are negligible for pions in asymmetric nuclear matter, are important as they have only one light quark. However, because of the large explicit symmetry breaking due to the finite strange quark mass higher order corrections in the chiral expansion are non-negligible, leading to an appreciable attractive scalar interaction for both the kaon and the antikaon in the medium. Available experimental data on both subthreshold kaon production and kaon flow are consistent with the presence of this attractive scalar interaction.

For etas, their properties are more related to the explicit breaking of the  $U_A(1)$  axial symmetry in QCD. If the  $U_A(1)$  symmetry is restored in medium as indicated by the instanton liquid model, then their masses are expected to decrease as well. Heavy ion experiments at the CERN-SPS, however, rule out a vast enhancement of the final-state eta yield.

In the case of vector mesons, we have discussed the rho, omega, and phi. At finite temperature, all model calculations agree with the current algebra result that the mass of the rho does not change to order  $T^2$ . To higher order in the temperature and at finite density QCD sum rules predict a dropping of the masses in the medium. Predictions from chiral models, on the other hand, tend to predict an increase of the rho mass with temperature. However, presently available dilepton data from CERN SPS heavy ion experiments are best described assuming a dropping rho mass in dense matter but the errors in the present exper-

imental data are too large to definitely exclude some more conventional explanations.

For baryons, particularly the nucleon, the change of their properties in the nuclear medium has been well-known in studies based on conventional many-body theory. On the other hand, the nucleon mass is also found to be reduced in the medium as a result of the scalar attractive potential related to the quark condensate. A clear separation of the vacuum effects due to the condensate from those of many-body interactions is a topic of current interest and has not yet been resolved. A dropping nucleon effective mass in medium seems to be required to explain the large enhancement of antiproton production in heavy ion collisions at subthreshold energies.

In principle lattice QCD calculations could help answer quite a few of these questions, although they are restricted to systems at finite temperature and vanishing baryon density. However, with the presently available computing power reliable quantitative predictions about in-medium properties of hadrons are still not available.

The study of in-medium properties of hadrons is a very active field, both theoretically and experimentally. While many questions are still open and require additional measurements and more careful calculations, it is undisputed that the study of in medium properties of hadrons provides us with an unique opportunity to further our understanding about the long range, nonperturbative aspects of QCD.

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## References

- [1] DeTar CE. 1995. *Quark Gluon Plasma 2* ed. R. Hwa, Singapore:

World Scientific

- [2] Drukarev EG, Levin EM. 1990. *Nucl. Phys. A* 511:679-700
- [3] Cohen TD, Furnstahl RJ, Griegel DK. 1992. *Phys. Rev. C* 45:1881-93
- [4] Li GQ, Ko CM. 1994. *Phys. Lett. B* 338:118-22
- [5] Brockmann R, Weise W. 1996. *Phys. Lett. B* 367:40-44.
- [6] Schäfer T, Shuryak EV. 1996. Instantons in QCD. *hep-ph/9610451*, to appear in *Rev. Mod. Phys*
- [7] Rho M. 1994. *Phys. Repts.* 240:1-142
- [8] Brown GE, Rho M. 1991. *Phys. Rev. Lett.* 66:2720-3
- [9] Aichelin J, Ko CM. 1985. *Phys. Rev. Lett.* 55:2661-4
- [10] Batko G, Cassing W, Mosel U, Niita K, Wolf G. 1991. *Phys. Lett. B* 256:331-6
- [11] Jeukenne JP, Lejeune A, and Mahaux C. 1976 *Phys. Rep.* 25:83-174
- [12] Mahaux C, Bortignon PF, Broglia RA, Dasso CH. 1985. *Phys. Rep.* 120:1-274
- [13] Brown GE, Weise W, Baym G, Speth J. 1987. *Comments Nucl. Part. Phys.* 17:39-62
- [14] Gelmini G, Ritzi B. 1995. *Phys. Lett. B* 357:431-4
- [15] Brown GE, Rho M. 1996. *Nucl. Phys. A* 596:503-14
- [16] Furnstahl RJ, Tang HB, Serot BD. 1995. *Phys. Rev. C* 52:1368-79
- [17] Stöcker H, Greiner W. 1986. *Phys. Rep.* 137:277-392
- [18] Bertsch GF, Das Gupta S. 1988. *Phys. Rep.* 160:189-233
- [19] Cassing W, Metag V, Mosel U, Niita K. 1990. *Phys. Rep.* 188:363-449
- [20] Aichelin J. 1991. *Phys. Rep.* 202:235-360
- [21] Blättel B., Koch V., and Mosel U. 1993. *Rep. Prog. Phys.* 56:1-62
- [22] Bonasera A, Gulminelli G, Molitoris J. 1994. *Phys. Rep.* 243:1-124
- [23] Ko CM, Li GQ. 1996. *J. Phys. G* 22:1673-725
- [24] Effenberger M, Hombach A, Teis S, Mosel M. 1996. *nucl-th/9607005*
- [25] Effenberger M, Hombach A, Teis S, Mosel M. 1996. *nucl-th/9610022*
- [26] Li BA, Bauer W, Ko CM. *Phys. Lett. B* 382:337-42
- [27] Matsui T, Satz H. 1986. *Phys. Lett. B* 178:416-422
- [28] Gerschel C, Hüfner J. 1992. *Z. Phys. C* 56:171-4
- [29] Gavin S, Vogt R. 1990. *Nucl. Phys. B* 345:104-24
- [30] Cassing W, Ko CM. 1996. *Phys. Lett. B*. In press.



- [31] 1996. *Proc. Quark Matter 96, Nucl. Phys. A610*, ed. P. Braun-Munzinger *et al.*
- [32] Kharzeev D, Satz H. 1995. In *Quark Gluon Plasma 2*, ed. R. Hwa, Singapore: World Scientific.
- [33] Ericson T, Weise W. 1988. *Pions and Nuclei*. Oxford: Clarendon Press
- [34] Delorme J, Ericson M, Ericson TEO. 1992. *Phys. Lett. B* 291:379-84
- [35] Thorsson V, Wirzba A. 1995. *Nucl. Phys. A* 589:633-48
- [36] Lutz M, Klimt S., Weise W. 1992. *Nucl. Phys. A* 542 521-58
- [37] Weinberg S. 1966. *Phys. Rev. Lett.* 17:616-21
- [38] Bertsch GF, Brown GE, Koch V, Li BA. 1988. *Nucl. Phys. A* 490:745-55
- [39] Xiong L, Ko CM, Koch V. 1993. *Phys. Rev. C* 47:788-94
- [40] Friedman B, Pandharipande VR, Usmani QN. 1981. *Nucl. Phys. A*, 372:483-95
- [41] Ericson M, Ericson TEO. 1966. *Ann. Phys. (N.Y.)* 36:323-362
- [42] Brown GE, Oset E, Vacas MV, Weise W. 1989. *Nucl. Phys. A* 505:823-34
- [43] Ko CM, Xia LH, Siemens PJ. 1989. *Phys. Lett. B* 231:16-20
- [44] Xia LH, Siemens PJ, Soyeur M. 1994. *Nucl. Phys. A* 578:493-510
- [45] Henning PA, Umezawa H. 1994. *Nucl. Phys. A* 571:617-644
- [46] Korpa CL, Malfliet R. 1995. *Phys. Rev. C* 52:2756-61
- [47] Helgeson J, Randrup J. 1995. *Ann. Phys. (NY)* 244:12-66
- [48] Hennio T, et al. 1993. *Phys. Lett. B* 303:236-39
- [49] C. Garde, private communication.
- [50] Gale C, Kapusta J. 1987. *Phys. Rev. C* 35:2107-2116
- [51] Korpa CL, Pratt S. 1990. *Phys. Rev. Lett.* 64:1502-5
- [52] Korpa CL, Xiong L, Ko CM, Siemens PJ. 1990. *Phys. Lett. B* 246:333-6
- [53] Roche G, *et al.* 1988. *Phys. Rev. Lett.* 61:1069-72
- [54] Naudet C, *et al.* 1988. *Phys. Rev. Lett.* 62:2652-55
- [55] Ehehalt W, Cassing W, Engel A, Mosel U, Wolf G. 1993. *Phys. Lett. B* 298:31-5
- [56] Zipprich J, Fuchs C, Lehmann E, Sehn L, Huang S W, Faessler A. 1997. *J. Phys. G* 23:L1-L6
- [57] Odyniec G. 1987. *Proc. 8th High Energy Heavy Ion Study*, ed. JW Harris, GJ Wozniak, pp. 215. Berkeley
- [58] Berg FD, *et al.* 1990. *Z.Phys. A* 340:297-302
- [59] Gerber P, Leutwyler H. 1989. *Nucl. Phys. B* 321:387-429
- [60] Shuryak E. 1990. *Phys. Rev. D* 42:1764-76
- [61] Shuryak E. 1991. *Nucl. Phys. A* 533:761-788

- [62] Koch V, Bertsch GF. 1993. *Nucl. Phys. A* 552:591-604
- [63] Song C. 1994. *Phys. Rev. D* 49:1556-65
- [64] Song C, Koch V, Lee SH, Ko CM. 1996. *Phys. Lett. B* 366:379-84
- [65] Koch V, Song C. 1996. *Phys. Rev. C* 54:1903-13
- [66] Song C, Koch V. 1996. *Phys. Rev. C* 54:3218-31
- [67] Stroebele H, *et al.* (NA35 collaboration). 1988. *Z.Phys. C* 38:89-96
- [68] Stroebele H, *et al.* (NA35 collaboration). 1990. *Nucl. Phys. A* 525:59c-66c
- [69] Kaplan DB, Nelson AE. 1986. *Phys. Lett. B* 175:57-63
- [70] Nelson AE, Kaplan DB. 1987. *Phys. Lett. B* 192:193-7
- [71] Politzer HD, Wise MB. 1991. *Phys. Lett. B* 273:156-62
- [72] Brown GE, Lee CH, Rho M, Thorsson V. 1994. *Nucl. Phys. A* 567:937-56
- [73] Jenkins E. 1992. *Nucl. Phys. B* 368:190-203
- [74] Martin, A.D. 1981. *Nucl. Phys. B* 179:33-48
- [75] Dalitz RH, Wong TC, Rajasekaran G. 1967. *Phys. Rev* 153:1617-23
- [76] Lee CH, Brown GE, Rho M. 1994. *Phys. Lett. B* 335:266-72
- [77] Kaiser N, Siegel PB, Weise W. 1995. *Nucl. Phys. A* 594:325-45
- [78] Koch V. 1994. *Phys. Lett. B* 337:7-11
- [79] Friedman E, Gal A, Batty CJ. 1993. *Phys. Lett. B* 308:6-10
- [80] Waas T, Kaiser N, Weise W. 1996. *Phys. Lett. B* 365:12-6
- [81] Brown GE, Thorsson V, Kubodera K, Rho M. 1992. *Phys. Lett. B* 291:355-62
- [82] Brown GE, Bethe HA. 1994. *Astrophys. Jour.* 423:659-64
- [83] Fang XS, Ko CM, Li GQ, Zheng YM. 1994. *Phys. Rev. C* 49:R608-11
- [84] Fang XS, Ko CM, Li GQ, Zheng YM. 1994. *Nucl. Phys. A* 575:766-90
- [85] Li GQ, Ko CM. 1995. *Phys. Lett. B* 349:405-10
- [86] Li GQ, Ko CM, Li BA. 1995. *Phys. Rev. Lett.* 74:235-8
- [87] Li GQ, Ko CM. 1995. *Nucl. Phys. A* 594:460-82
- [88] Miskowiec D *et al.*, *Phys. Rev. Lett.* 72:3650-3
- [89] Ritman J, the FOPI collaboration. 1995. *Z.Phys. A* 352:355-7
- [90] David C, Hartnack C, Kerveno M, Le Pallec JC, Aichelin J. 1996. *Nucl. Phys. A*. Submitted
- [91] Maruyama T, Cassing W, Mosel U, Teis S, Weber K. 1994. *Nucl. Phys. A* 573:653-75
- [92] Gibson BF, Hungerford EV. 1995 *Phys. Rep.* 257:349
- [93] Chen CM, Ernst DJ. 1992. *Phys. Rev. C* 45:2019-22

- [94] Brown GE, Dover CB, Siegel PB, Weise W. 1988. *Phys. Rev. Lett.* 60:2723-26
- [95] Li GQ, Ko CM, Fang XS. 1994. *Phys. Lett. B* 329:149-56
- [96] Li GQ, Ko CM. 1996. *Phys. Rev. C* 54:R2159-R2162
- [97] Schröter A, *et al.* 1994. *Z.Phys. A* 350:101-13
- [98] Ritman. 1996. private communications
- [99] Shuryak E. 1994. *Comm. Nucl. Part. Phys.* 21:235-248
- [100] Schäfer T. 1996. *Phys. Lett. B* 389:455-51
- [101] Kapusta J, Kharzeev D, McLerran L. 1996. *Phys. Rev. D* 53:5028-33
- [102] Huang Z, Wang XN. 1996. *Phys. Rev. D* 53:5034-41
- [103] Berg FD, *et al.* 1994. *Phys. Rev. Lett.* 72:977-80
- [104] Drees A. 1996. *Phys. Lett. B* 388:380-3
- [105] Koch V. 1990. *Proc. Pittsburgh Workshop on Soft Lepton Pair and Photon Production*, ed. JA Thompson, p. 251. New York: Nova
- [106] Wolf G, Cassing W, Mosel U. 1993. *Nucl. Phys. A* 552:549-70
- [107] Li GQ, Ko CM. 1995. *Nucl. Phys. A* 583:731-48
- [108] Weinberg S. 1967. *Phys. Rev. Lett.* 18:507-9
- [109] Asakawa M, Ko CM. 1994. *Nucl. Phys. A* 572:732-48
- [110] Song C. 1996. *Phys. Lett. B* 388:141-6
- [111] Adami C, Brown GE. 1993. *Phys. Rep.* 234:1-71
- [112] Friman B, Rho M. 1996. *Nucl. Phys. A* 606:303-19
- [113] Brown GE, Buballa M, Rho M. 1996. *Nucl. Phys. A* 609:519-36
- [114] Donoghue J, Golowich E, Holstein B. 1992. *Dynamics of the Standard Model*. Cambridge: Cambridge University Press, Cambridge, U.K.
- [115] Pisarski R. 1982. *Phys. Lett. B* 110:155-59
- [116] Koch V, Brown GE. 1993. *Nucl. Phys. A* 560:345-64
- [117] Sakurai JJ. 1969. *Currents and mesons*. Chicago: University of Chicago Press, Chicago
- [118] Gale C, Kapusta JJ. 1991. *Nucl. Phys. B* 357:65-89
- [119] Reinders LJ, Rubinstein H, Yazaki S. 1985. *Phys. Rep.* 127:1-97
- [120] Kapusta JJ, Shuryak EV. 1994. *Phys. Rev. D* 49:4694-704
- [121] Dey M, Eletzky VL, Ioffe BL. 1990. *Phys. Lett. B* 252:620-4
- [122] Eletsky VL, Ioffe BL. 1995. *Phys. Rev. D* 51:2371-6
- [123] Pisarski RD. 1995. *Phys. Rev. D* 52:R3773-R3776
- [124] Song C. 1993. *Phys. Rev. D* 48:1375-89
- [125] Song CS, Xia PW, Ko CM. 1995. *Phys. Rev. C* 52:408-11
- [126] Rapp R, Wambach J. 1993. *Phys. Lett. B* 315:220-5
- [127] Chanfray G, Schuck P. 1993. *Nucl. Phys. A* 555:329-54

- [128] Herrmann M, Friman B, Nörenberg W. 1993. *Nucl. Phys. A* 560:411-36
- [129] Asakawa M, Ko CM, Levai P, Qiu XJ. *Phys. Rev. C* 46:R1159-R1162
- [130] Klingl F, Weise W. 1996. *Nucl. Phys. A* 606:329-38
- [131] Klingl F, Weise W. 1997. Proceedings of the Intl. Workshop XXV on Gross Properties of Nuclei and Nuclear Excitations ed. H. Feldmeier and W. Nörenberg, Hirschegg, Austria, 1997.
- [132] Klingl F, Weise W. in preparation
- [133] Hatsuda T, Lee SH. 1992. *Phys. Rev. C* 46:R34-R38
- [134] Asakawa M, Ko CM. 1993. *Nucl. Phys. A* 560:399-410
- [135] Friman B, Soyeur M. 1996. *Nucl. Phys. A* 600:477-90
- [136] Eletsky VL, Ioffe BL. 1996. *hep-ph/9609229*
- [137] Friman B, Pirner H. 1997 *nucl-th/9701016*
- [138] Xiong L, Wu ZG, Ko CM, Wu JQ. 1990. *Nucl. Phys. A* 512:772-86
- [139] Wolf G, Batko G, Cassing W, Mosel U, Niita K, Schäfer M. 1990. *Nucl. Phys. A* 517:615-38
- [140] Porter R.J. et al. 1997, LBNL-Report LBNL-39957, submitted to *Phys. Rev. Lett.*
- [141] Drees A. 1997 Proc. Quark Matter 96, Nucl. Phys. A610, ed. P. Braun-Munzinger
- [142] Agakichiev G, *et al.* 1995. *Phys. Rev. Lett.* 75:1272-5
- [143] Masera M for the HELIOS Collaboration. 1995. *Nucl. Phys. A* 590:93c-102c
- [144] Cassing W, Ehhalt W, Ko CM. 1995. *Phys. Lett. B* 363:35-40
- [145] Srivastava DK, Sinha B, Gale C. *Phys. Rev. C* 53:R567-R571
- [146] Haglin K. 1996. *Phys. Rev. C* 53:R2606-R2609
- [147] Baier R, Dirks M, Redlich R. 1996. *hep-ph/9610210*
- [148] Chanfray G, Rapp R, Wambach J. 1996. *Phys. Rev. Lett.* 76:368-71
- [149] LI GQ, Ko CM, Brown GE. 1995. *Phys. Rev. Lett.* 75:4007-10
- [150] LI GQ, Ko CM, Brown GE. 1996. *Nucl. Phys. A* 606:568-606
- [151] LI GQ, Ko CM, Brown GE, Sorge H. 1996. *Nucl. Phys. A* 611:539-67
- [152] Song CS, Lee SH, Ko CM. 1995. *Phys. Rev. C* 52:R476-9
- [153] Kharzeev, Satz H. 1994. *Phys. Lett. B* 340:167-70
- [154] Freedom BM. 1995. *Proc. Internal. Workshop XXXIII on Gross Properties of Nuclei and Nuclear Excitations*, pp. 273-82. GSI
- [155] Stroth J. 1995. *Proc. Internal. Workshop XXXIII on Gross Properties of Nuclei and Nuclear Excitations*, pp. 202-17. GSI

- [156] Jean HC, Piekarewicz J, Williams AG. 1994. *Phys. Rev. C* 49:1981-8
- [157] Shiomo H, Hatsuda T. 1994. *Phys. Lett. B* 334:281-6
- [158] Ko CM, Lévai P, Qiu XJ. 1992. *Phys. Rev. C* 45:1400-2
- [159] Kuwabara H, Hatsuda T. 1995. *Prog. Theo. Phys.* 94:1163-7
- [160] Akiba Y, *et al.*. 1996. *Phys. Rev. Lett.* 76:2021-24
- [161] Asakawa M, Ko CM. 1994. *Phys. Lett. B* 322:33-7
- [162] Asakawa M, Ko CM. 1994. *Phys. Rev. C* 50:3064-8
- [163] Cohen TD, Furnstahl RJ, Griegel DK. 1991. *Phys. Rev. Lett.* 67:961-4
- [164] Jin XM, Cohen TD, Furnstahl RJ, Griegel DK. 1993. *Phys. Rev. C* 47:2882-900
- [165] Jin XM, Nielsen M, Cohen TD, Furnstahl RJ, Griegel DK. 1994. *Phys. Rev. C* 49:464-77
- [166] Hatsuda T, Higaasen T, Prakash M. 1990. *Phys. Rev. C* 42:2212-21
- [167] Adami C, Brown GE. 1990 *Z. Phys. A* 340:93-100
- [168] Serot BD, Walecka JD. 1986. *Adv. Nucl. Phys.* 16:1
- [169] Machleidt R. 1989. *Adv. Nucl. Phys.* 19:189-376
- [170] Ray L, Hoffmann GW, Coker WR. 1992. *Phys. Rep.* 212:223-328
- [171] Arnold LG, Clark BC, Mercer RL. 1979. *Phys. Rev. C* 19:917-22
- [172] Kobbs AM, Cooper ED, Johansson JI, Sherif HS. 1985. *Nucl. Phys. A* 445:605-24
- [173] Hama S, Clark BC, Cooper ED, Sherif HS, Mercer RL. 1990. *Phys. Rev. C* 41:2737-55
- [174] Cooper ED, Hama S, Clark BC, Mercer RL. 1993. *Phys. Rev. C* 47:297-311
- [175] Li GQ, Ko CM. 1994. *Phys. Rev. C* 50:1725-8
- [176] Li GQ, Ko CM, Fang XS, Zheng YM. 1994. *Phys. Rev. C* 49:1139-48
- [177] Teis S, Cassing W, Maruyama T, Mosel U. 1994. *Phys. Rev. C* 50:388-405
- [178] Batko G, Faessler A, Huang SW, Lehmann E, Puri RK. 1994. *J. Phys. G* 20:461-5
- [179] Schröter A, *et al.* 1994. *Nucl. Phys. A* 553:775c-778c
- [180] Baldin AA, *et al.* 1992. *Nucl. Phys. A* 519:407c-411c
- [181] Koch V, Ko CM, Brown GE. 1991. *Phys. Lett. B* 265:29-34
- [182] Spieles C, Bleicher M, Jahns A, Mattiello R, Sorge H, Stöcker H, Greiner W. 1996. *Phys. Rev. C* 53:2011-3

- [183] Sexton J, Vaccarino A, Weingarten D, Butler F, Chen H. 1994. *Nucl. Phys. B* 430:179-228
- [184] Eletskii VL, Ioffe BL. 1988. *Sov. J. Nucl. Phys* 48:602-
- [185] DeTar CE, Kogut J. 1987. *Phys. Rev. D* 36:2828-39
- [186] Born KD, Gupta S, Irbäck A, Karsch F, Laermann E, Petersson B, Satz H, (MTC Collaboration). 1991. *Phys. Rev. Lett.* 67:302-5
- [187] Boyd G, Gupta S, Karsch F, Laermann E, Peterson B, Redlich K. 1995. *Phys. Lett. B* 349:170-6
- [188] Learmann E, 1996 *Proc. Quark Matter 96, Heidelberg, to appear in Nucl. Phys. A*, ed. J Specht. North Holland